

# Minimal $\mathcal{N} = 4$ topologically massive supergravity

Sergei M. Kuzenko<sup>a</sup>, Joseph Novak<sup>b</sup> and Ivo Sachs<sup>c</sup>

<sup>a</sup>*School of Physics M013, The University of Western Australia  
35 Stirling Highway, Crawley W.A. 6009, Australia*

<sup>b</sup>*Max-Planck-Institut für Gravitationsphysik, Albert-Einstein-Institut  
Am Mühlenberg 1, D-14476 Golm, Germany*

<sup>c</sup>*Arnold Sommerfeld Center for Theoretical Physics, Ludwig-Maximilians-Universität  
Theresienstrae 37, D-80333 München, Germany*

sergei.kuzenko@uwa.edu.au, joseph.novak@aei.mpg.de,  
ivo.sachs@physik.uni-muenchen.de

## Abstract

Using the superconformal framework, we construct a new off-shell model for topologically massive  $\mathcal{N} = 4$  supergravity which is minimal in the sense that it makes use of a single compensating vector multiplet. Our theory provides a counterexample to the common lore that two compensating multiplets are required within the conformal approach to supergravity with eight supercharges in diverse dimensions. All solutions in this theory correspond to non-conformally flat superspaces. Its maximally supersymmetric solutions include the so-called critical (4,0) anti-de Sitter superspace introduced in arXiv:1205.4622. Other maximally supersymmetric solutions describe warped critical (4,0) anti-de Sitter superspaces. We also propose a dual formulation for the theory in which the vector multiplet is replaced with an off-shell hypermultiplet. Upon elimination of the auxiliary fields belonging to the hypermultiplet and imposing certain gauge conditions, the dual action reduces to the one introduced in arXiv:1605.00103.

# 1 Introduction

A unique feature of three spacetime dimensions (3D) is the existence of topologically massive Yang-Mills and gravity theories. They are obtained by augmenting the usual Yang-Mills action or the gravitational action by a gauge-invariant topological mass term. Such a mass term coincides with a non-Abelian Chern-Simons action in the Yang-Mills case [1, 2, 3, 4] and with a Lorentzian Chern-Simons term in the case of gravity [3, 4]. Without adding the Lorentzian Chern-Simons term, the pure gravity action propagates no local degrees of freedom. The Lorentzian Chern-Simons term can be interpreted as the action for conformal gravity in three dimensions [3, 5, 6].<sup>1</sup>

Topologically massive theories of gravity possess supersymmetric extensions. In particular, topologically massive  $\mathcal{N} = 1$  supergravity was introduced in [9] and its cosmological extension followed in [10]. The off-shell formulations for topologically massive  $\mathcal{N}$ -extended supergravity theories were presented in [11] for  $\mathcal{N} = 2$  and in [12] for  $\mathcal{N} = 3$  and  $\mathcal{N} = 4$ . In all of these theories, the action functional is a sum of two terms, one of which is the action for pure  $\mathcal{N}$ -extended supergravity (Poincaré or anti-de Sitter) and the other is the action for  $\mathcal{N}$ -extended conformal supergravity. The off-shell actions for  $\mathcal{N}$ -extended supergravity theories in three dimensions were given in [13] for  $\mathcal{N} = 1$ , [14, 15] for  $\mathcal{N} = 2$ , and [14] for the cases  $\mathcal{N} = 3, 4$ . The off-shell actions for  $\mathcal{N}$ -extended conformal supergravity were given in [5] for  $\mathcal{N} = 1$ , [16] for  $\mathcal{N} = 2$ , and [17] for the cases  $\mathcal{N} = 3, 4$ . The latter work made use of the formulation for  $\mathcal{N}$ -extended conformal supergravity presented in [18].

The off-shell structure of 3D  $\mathcal{N} = 4$  supergravity [14] is analogous to that of 4D  $\mathcal{N} = 2$  supergravity (see, e.g., [19] for a pedagogical review) in the sense that two superconformal compensators are required (for instance, two off-shell vector multiplets, one of which is self-dual and the other anti-self-dual) in order to realise pure Poincaré or anti-de Sitter (AdS) supergravity theories. We recall that the equations of motion for pure  $\mathcal{N} = 4$  Poincaré or AdS supergravity are inconsistent if one makes use of a single compensator [12]. By construction, the off-shell topologically massive  $\mathcal{N} = 4$  supergravity theory of [12] makes use of two compensators. However, in [20] the consistent system of dynamical equations was proposed for topologically massive  $\mathcal{N} = 4$  AdS supergravity with a single compensating hypermultiplet, following earlier work in [21, 22, 23] on ABJ(M) models. A peculiar feature of this model, like those considered in [21, 22, 23], is that it has no free parameter. Consequently the dimensionless combination,  $\mu\ell$ , of mass  $\mu$  and AdS radius  $\ell$  takes a fixed value,  $\mu\ell = 1$ ,

---

<sup>1</sup>The usual Einstein-Hilbert action for 3D gravity with a cosmological term can also be interpreted as the Chern-Simons action for the anti-de Sitter group [7, 8].

as in chiral gravity [24]. In [20] a supergravity action functional was also postulated to generate the dynamical equations given. This action was claimed to be off-shell without giving technical details. In this paper we propose a new off-shell model for topologically massive  $\mathcal{N} = 4$  supergravity which is minimal in the sense that it makes use of a single compensating vector multiplet. The theory is consistent only if the term corresponding to  $\mathcal{N} = 4$  conformal supergravity is turned on. An important maximally supersymmetric solution for this theory is the so-called critical (4,0) AdS superspace introduced in [25]. Our supergravity theory is first presented in a manifestly supersymmetric form, and then its action functional is reduced to components. By choosing appropriate gauge conditions at the component level and performing a duality transformation, we show how to reduce our off-shell supergravity action to the one postulated in [20].

This paper is organised as follows. In section 2 we recall the superspace geometry of the  $\mathcal{N} = 4$  vector multiplets and the corresponding locally supersymmetric actions. In section 3 we present two models for minimal topologically massive  $\mathcal{N} = 4$  supergravity, analyse their equations of motion and give a brief discussion of the maximally supersymmetric solutions. Section 4 is devoted to the component structure of minimal topologically massive  $\mathcal{N} = 4$  supergravity. Concluding comments are given in section 5. The main body of the paper is accompanied with three technical appendices. The essential details of the known superspace formulations for  $\mathcal{N} = 4$  conformal supergravity are collected in Appendices A and B. Some useful super-Weyl gauge conditions in  $\text{SO}(4)$  superspace and their implications are given in Appendix C.

## 2 The $\mathcal{N} = 4$ vector multiplets

There are two inequivalent irreducible  $\mathcal{N} = 4$  vector multiplets in three dimensions, self-dual and anti-self-dual ones, as discovered by Brooks and Gates [26]. In this section we review the superspace geometry of these supermultiplets in the presence of  $\mathcal{N} = 4$  conformal supergravity [14, 18] and the corresponding locally supersymmetric actions [14].

Throughout this paper we make use of both the  $\text{SO}(4)$  superspace formulation of conformal supergravity, which was sketched in [27] and fully developed in [14], and the conformal superspace formulation presented in [18]. These formulations are related to each other since  $\text{SO}(4)$  superspace may be viewed as a gauge fixed version of the  $\mathcal{N} = 4$  conformal superspace [18]. Due to this reason, we will first start by formulating

vector multiplets in conformal superspace. We refer the reader to Appendix A for the salient details of the conformal superspace formulation. The geometry of SO(4) superspace is briefly reviewed in Appendix B.

## 2.1 Kinematics

To describe an Abelian vector multiplet in a curved superspace  $\mathcal{M}^{3|8}$  parametrised by coordinates  $z^M = (x^m, \theta^\mu_x)$ , we introduce gauge covariant derivatives

$$\nabla = E^A \nabla_A, \quad \nabla_A = (\nabla_a, \nabla_\alpha^I) := \nabla_A - V_A \mathbf{Z}, \quad [\mathbf{Z}, \nabla_A] = 0, \quad (2.1)$$

with  $E^A = dZ^M E_M^A$  the superspace vielbein,  $\nabla_A$  the superspace covariant derivatives (A.2) obeying the (anti-)commutation relations (A.4), and  $V = E^A V_A$  the gauge connection associated with  $\mathbf{Z}$ . The gauge transformation of  $V$  is

$$\delta V = d\tau, \quad (2.2)$$

where the gauge parameter  $\tau(z)$  is an arbitrary scalar superfield.

The algebra of gauge covariant derivatives is

$$\begin{aligned} [\nabla_A, \nabla_B] = & -T_{AB}^C \nabla_C - \frac{1}{2} R(M)_{AB}{}^{cd} M_{cd} - \frac{1}{2} R(N)_{AB}{}^{PQ} N_{PQ} - R(\mathbb{D})_{AB} \mathbb{D} \\ & - R(S)_{AB}{}^\gamma S_\gamma^I - R(K)_{AB}{}^c K_c - F_{AB} \mathbf{Z}, \end{aligned} \quad (2.3)$$

where the torsion and curvatures are those of conformal superspace but with  $F_{AB}$  corresponding to the gauge covariant field strength  $F = \frac{1}{2} E^B \wedge E^A F_{AB} = dV$ . The field strength  $F_{AB}$  satisfies the Bianchi identity

$$dF = 0, \quad \nabla_{[A} F_{BC]} + T_{[AB}{}^D F_{D|C]} = 0 \quad (2.4)$$

and must be subject to covariant constraints to describe an irreducible vector multiplet.

In order to describe an  $\mathcal{N} = 4$  vector multiplet, the superform  $F$  is subject to the constraint (see [14] for more details)

$$F_{\alpha\beta}^{IJ} = -2i\varepsilon_{\alpha\beta} G^{IJ}, \quad G^{IJ} = -G^{JI}, \quad (2.5a)$$

and then the Bianchi identity fixes the remaining components of  $F$  to be

$$F_{a\beta}^J = \frac{1}{3} (\gamma_a)_\beta{}^\gamma \nabla_{\gamma K} G^{JK}, \quad (2.5b)$$

$$F_{ab} = -\frac{i}{48}\varepsilon_{abc}(\gamma^c)^{\alpha\beta}[\nabla_\alpha^K, \nabla_\beta^L]G_{KL} , \quad (2.5c)$$

where  $G^{IJ}$  is primary and of dimension 1,

$$S_\alpha^I G^{JK} = 0 , \quad K_a G^{IJ} = 0 , \quad \mathbb{D}G^{IJ} = G^{IJ} . \quad (2.6)$$

Moreover, the field strength  $G^{IJ}$  is constrained by the dimension-3/2 Bianchi identity

$$\nabla_\gamma^I G^{JK} = \nabla_\gamma^{[I} G^{JK]} - \frac{2}{3}\delta^{I[J} \nabla_{\gamma L} G^{K]L} . \quad (2.7)$$

It is well known (see [14] and references therein) that the constraint (2.7) defines a reducible off-shell supermultiplet.<sup>2</sup> The point is that the Hodge-dual of  $G^{IJ}$ ,

$$\tilde{G}^{IJ} := \frac{1}{2}\varepsilon^{IJKL}G_{KL} , \quad (2.8)$$

obeys the same constraint as  $G^{IJ}$  does,

$$\nabla_\gamma^I \tilde{G}^{JK} = \nabla_\gamma^{[I} \tilde{G}^{JK]} - \frac{2}{3}\delta^{I[J} \nabla_{\gamma L} \tilde{G}^{K]L} , \quad (2.9a)$$

where  $\varepsilon^{IJKL}$  is the Levi-Civita tensor. As a result one may constrain the field strength  $G^{IJ}$  to be self-dual,  $\tilde{G}^{IJ} = G^{IJ}$  or anti-self-dual,  $\tilde{G}^{IJ} = -G^{IJ}$ . These choices correspond to two different irreducible off-shell  $\mathcal{N} = 4$  vector multiplets, which we denote by  $G_+^{IJ}$  and  $G_-^{IJ}$ , respectively. In what follows we will make use of an (anti-)self-dual Abelian vector multiplet such that its field strength  $G_\pm^{IJ}$  is nowhere vanishing,  $G_\pm^2 := \frac{1}{2}G_\pm^{IJ}G_{\pm IJ} \neq 0$ .

When working with  $\mathcal{N} = 4$  supersymmetric theories, a powerful technical tool is the isospinor notation based on the isomorphism  $\text{SO}(4) \cong (\text{SU}(2)_L \times \text{SU}(2)_R)/\mathbb{Z}_2$ , which allows one to replace each  $\text{SO}(4)$  vector index with a pair of isospinor ones. In defining the isospinor notation, we follow [14] and associate with a real  $\text{SO}(4)$  vector  $V_I$  a second-rank isospinor  $V_{i\bar{i}}$  defined as

$$V_I \rightarrow V_{i\bar{i}} := (\tau^I)_{i\bar{i}} V_I , \quad V_I = \tau_I^{i\bar{i}} V_{i\bar{i}} , \quad (V_{i\bar{i}})^* = V^{i\bar{i}} , \quad (2.10)$$

where we have introduced the  $\tau$ -matrices

$$(\tau^I)_{i\bar{i}} = (\mathbb{1}, i\sigma_1, i\sigma_2, i\sigma_3) , \quad I = 1, \dots, 4 , \quad i = 1, 2 , \quad \bar{i} = \bar{1}, \bar{2} . \quad (2.11)$$

The isospinor indices of  $\text{SU}(2)_L$  and  $\text{SU}(2)_R$  spinors  $\psi_i$  and  $\chi_{\bar{i}}$ , respectively, are raised and lowered using the antisymmetric tensors  $\varepsilon^{ij}, \varepsilon_{ij}$  and  $\varepsilon^{\bar{i}\bar{j}}, \varepsilon_{\bar{i}\bar{j}}$  (normalised by  $\varepsilon^{12} = \varepsilon_{21} = \varepsilon^{\bar{1}\bar{2}} = \varepsilon_{\bar{2}\bar{1}} = 1$ ) according to

$$\psi^i = \varepsilon^{ij}\psi_j , \quad \psi_i = \varepsilon_{ij}\psi^j , \quad \chi^{\bar{i}} = \varepsilon^{\bar{i}\bar{j}}\chi_{\bar{j}} , \quad \chi_{\bar{i}} = \varepsilon_{\bar{i}\bar{j}}\chi^{\bar{j}} . \quad (2.12)$$

---

<sup>2</sup>Such a long  $\mathcal{N} = 4$  supermultiplet naturally originates upon reduction of any off-shell  $\mathcal{N} > 4$  vector multiplet to  $\mathcal{N} = 4$  superspace [28].

We then have the following dictionary:

$$V^I U_I = V^{\bar{i}\bar{i}} U_{\bar{i}\bar{i}} , \quad (2.13a)$$

$$A_{\bar{i}\bar{j}\bar{j}} := A_{IJ}(\tau^I)_{\bar{i}\bar{i}}(\tau^J)_{\bar{j}\bar{j}} = \varepsilon_{ij} A_{\bar{i}\bar{j}} + \varepsilon_{\bar{i}\bar{j}} A_{ij} , \quad A_{ij} = A_{ji} , \quad A_{\bar{i}\bar{j}} = A_{\bar{j}\bar{i}} , \quad (2.13b)$$

$$\frac{1}{2} A^{IJ} B_{IJ} = A^{ij} B_{ij} + A^{\bar{i}\bar{j}} B_{\bar{i}\bar{j}} , \quad (2.13c)$$

$$\varepsilon_{\bar{i}\bar{j}\bar{j}k\bar{k}l\bar{l}} = \varepsilon_{ij}\varepsilon_{kl}\varepsilon_{\bar{i}\bar{l}}\varepsilon_{\bar{j}\bar{k}} - \varepsilon_{il}\varepsilon_{jk}\varepsilon_{\bar{i}\bar{j}}\varepsilon_{\bar{k}\bar{l}} , \quad (2.13d)$$

where  $V^I$  and  $U^I$  are  $\text{SO}(4)$  vectors,  $A^{IJ}$  and  $B^{IJ}$  are anti-symmetric second-rank  $\text{SO}(4)$  tensors. The left-hand side of (2.13d) is the Levi-Civita tensor in the isospinor notation.

In the isospinor notation, the self-dual  $(G_+^{IJ})$  and anti-self-dual  $(G_-^{IJ})$  vector multiplets take the form

$$G_+^{\bar{i}\bar{j}\bar{j}} = -\varepsilon^{ij} G^{\bar{i}\bar{j}} , \quad G_-^{\bar{i}\bar{j}\bar{j}} = -\varepsilon^{\bar{i}\bar{j}} G^{ij} , \quad (2.14)$$

and the Bianchi identity (2.7) turns into

$$\nabla_\alpha^{(\bar{i}\bar{i}} G^{kl)} = 0 , \quad \nabla_\alpha^{i(\bar{i}} G^{\bar{k}\bar{l})} = 0 . \quad (2.15)$$

At this stage it is useful to introduce left and right isospinor variables  $v_L := v^i \in \mathbb{C}^2 \setminus \{0\}$  and  $v_R := v^{\bar{i}} \in \mathbb{C}^2 \setminus \{0\}$ , which can be used to package the anti-self-dual field strength  $G^{ij}$  and the self-dual field strength  $G^{\bar{i}\bar{j}}$  into fields without isospinor indices,  $G_L^{(2)}(v_L) := G_{ij} v^i v^j$  and  $G_R^{(2)}(v_R) := G_{\bar{i}\bar{j}} v^{\bar{i}} v^{\bar{j}}$ , respectively. The same isospinor variables can be used to define two different subsets,  $\nabla_\alpha^{(1)\bar{i}}$  and  $\nabla_\alpha^{(\bar{1})i}$ , in the set of spinor covariant derivatives  $\nabla_\alpha^{\bar{i}\bar{i}}$  by the rule

$$\nabla_\alpha^{(1)\bar{i}} := v_i \nabla_\alpha^{\bar{i}\bar{i}} , \quad \nabla_\alpha^{(\bar{1})i} := v_{\bar{i}} \nabla_\alpha^{\bar{i}\bar{i}} . \quad (2.16)$$

It follows from (A.17) that the operators  $\nabla_\alpha^{(1)\bar{i}}$  obey the anti-commutation relations:

$$\begin{aligned} \{\nabla_\alpha^{(1)\bar{i}}, \nabla_\beta^{(1)\bar{j}}\} &= 2i\varepsilon_{\alpha\beta}\varepsilon^{\bar{i}\bar{j}} W L^{(2)} + i\varepsilon_{\alpha\beta}\varepsilon^{\bar{i}\bar{j}} \nabla^{\gamma(1)}_{\bar{k}} W S_\gamma^{(1)\bar{k}} \\ &\quad - \frac{1}{4}\varepsilon_{\alpha\beta}\varepsilon^{\bar{i}\bar{j}} \nabla_{\gamma^{(1)}\bar{k}} \nabla_\delta^{(1)\bar{k}} W K^{\gamma\delta} , \end{aligned} \quad (2.17)$$

where  $L^{(2)} = v_i v_j L^{ij}$  and  $S_\alpha^{(1)\bar{i}}$  is defined similarly to  $\nabla_\alpha^{(1)\bar{i}}$ . The rationale for the definitions given is that the constraints (2.15) now become the analyticity conditions

$$\nabla_\alpha^{(1)\bar{i}} G_L^{(2)} = 0 , \quad \nabla_\alpha^{(\bar{1})i} G_R^{(2)} = 0 . \quad (2.18)$$

which tell us that each of  $G_L^{(2)}$  and  $G_R^{(2)}$  depends on half the Grassmann coordinates. The constraints (2.18) do not change under re-scalings  $v^i \rightarrow c_L v^i$  and  $v^{\bar{i}} \rightarrow c_R v^{\bar{i}}$ ,

with  $c_L, c_R \in \mathbb{C} \setminus \{0\}$ , with respect to which  $G_L^{(2)}(v_L)$  and  $G_R^{(2)}(v_R)$  are homogeneous polynomials of degree two. We see that the isospinor variables  $v_L$  and  $v_R$  are defined modulo the equivalence relations  $v^i \sim c_L v^i$  and  $v^{\bar{i}} \sim c_R v^{\bar{i}}$ , and therefore they parametrise identical complex projective spaces  $\mathbb{CP}_L^1$  and  $\mathbb{CP}_R^1$ . The superfields  $G_L^{(2)}(v_L)$  and  $G_R^{(2)}(v_R)$  are naturally defined on curved  $\mathcal{N} = 4$  projective superspace  $\mathcal{M}^{3|8} \times \mathbb{CP}_L^1 \times \mathbb{CP}_R^1$  introduced in [14].

The field strengths  $G_L^{(2)}(v_L)$  and  $G_R^{(2)}(v_R)$  are examples of the covariant projective multiplets introduced in [14] in  $\text{SO}(4)$  superspace and later reformulated in [12] within the conformal superspace setting. There are two types of covariant projective multiplets, the left and right ones. A left projective multiplet of weight  $n$ ,  $Q_L^{(n)}(v_L)$ , is a superfield that is defined on some open domain of  $\mathbb{C}^2 \setminus \{0\}$  and possesses the following four properties. Firstly, it is a primary superfield,

$$S_{\alpha}^{\bar{i}i} Q_L^{(n)} = 0, \quad K_a Q_L^{(n)} = 0. \quad (2.19)$$

Secondly, it is subject to the constraint

$$\nabla_{\alpha}^{(1)\bar{i}} Q_L^{(n)} = 0. \quad (2.20)$$

Thirdly, it is a *holomorphic* function of  $v_L$ . Fourthly, it is *homogeneous* function of  $v_L$  of degree  $n$ ,

$$Q_L^{(n)}(c v_L) = c^n Q_L^{(n)}(v_L), \quad c \in \mathbb{C} \setminus \{0\}. \quad (2.21)$$

Every left projective multiplet is inert with respect to  $\text{SU}(2)_R$  and transforms under  $\text{SU}(2)_L$  as

$$\delta_{\Lambda} Q_L^{(n)} = \Lambda^{ij} L_{ij} Q_L^{(n)}, \quad (2.22a)$$

$$\Lambda^{ij} L_{ij} Q_L^{(n)} = -(\Lambda_L^{(2)} \partial_L^{(-2)} - n \Lambda_L^{(0)}) Q_L^{(n)}, \quad (2.22b)$$

where we have defined

$$\Lambda_L^{(2)} := \Lambda^{ij} v_i v_j, \quad \Lambda_L^{(0)} := \frac{v_i u_j}{(v_L, u_L)} \Lambda^{ij} \quad (2.23)$$

and made use of the differential operator

$$\partial_L^{(-2)} := \frac{1}{(v_L, u_L)} u^i \frac{\partial}{\partial v^i}, \quad (v_L, u_L) = v^i u_i. \quad (2.24)$$

Here we have also introduced a second left isospinor variable  $u_L := u^i$  which is restricted to be linearly independent of  $v_L$ , that is  $(v_L, u_L) \neq 0$ . One may see that  $L^{(2)} Q_L^{(n)} = 0$ , and therefore the integrability condition  $\{\nabla_{\alpha}^{(1)\bar{i}}, \nabla_{\beta}^{(1)\bar{j}}\} Q^{(n)} = 0$  for the constraint (2.20) holds, in accordance with (2.17). The right projective multiplets

are defined similarly. The covariant projective multiplets  $G_L^{(2)}(v_L)$  and  $G_R^{(2)}(v_R)$  are known as the left and right  $\mathcal{O}(2)$  multiplets, respectively.

As shown in [14] the self-dual vector multiplet,  $G_R^{(2)}(v_R)$ , can be described in terms of a gauge prepotential  $\mathcal{V}_L(v_L)$ , which is a left weight-0 *tropical multiplet* and is real with respect to the analyticity preserving conjugation called the smile conjugation. The interested reader is referred to [14] for the technical details. Similar properties hold for the anti-self-dual vector multiplet except all ‘left’ objects have to be replaced by ‘right’ ones and vice versa.

## 2.2 Dynamics

General off-shell matter couplings in  $\mathcal{N} = 4$  supergravity were constructed in [14]. The action for such a supergravity-matter system may be represented as a sum of two terms (one of which may be absent),

$$S = S_L + S_R . \quad (2.25)$$

The left  $S_L$  and right  $S_R$  actions, are naturally formulated in curved  $\mathcal{N} = 4$  projective superspace. The left action has the form

$$S_L = \frac{1}{2\pi} \oint (v_L, dv_L) \int d^{3|8}z E C_L^{(-4)} \mathcal{L}_L^{(2)} , \quad E^{-1} = \text{Ber}(E_A^M) , \quad (2.26)$$

where the Lagrangian  $\mathcal{L}_L^{(2)}(v_L)$  is a real left projective multiplet of weight 2, and  $d^{3|8}z$  denotes the full superspace integration measure,  $d^{3|8}z := d^3x d^8\theta$ . Furthermore, the *model-independent* primary isotwistor superfield  $C_L^{(-4)}(v_L)$  has dimension  $-2$ , i.e.  $\mathbb{D}C_L^{(-4)} = -2C_L^{(-4)}$ . It is defined to be real with respect to the smile-conjugation defined in [14] and obeys the differential equation

$$\Delta_L^{(4)} C_L^{(-4)} = 1 . \quad (2.27)$$

Here  $\Delta_L^{(4)}$  denotes the following fourth-order operator<sup>3</sup>

$$\Delta_L^{(4)} = \frac{1}{96} \left( \nabla^{(2)\bar{i}\bar{j}} \nabla_{\bar{i}\bar{j}}^{(2)} - \nabla^{(2)\alpha\beta} \nabla_{\alpha\beta}^{(2)} \right) = \frac{1}{48} \nabla^{(2)\bar{i}\bar{j}} \nabla_{\bar{i}\bar{j}}^{(2)} , \quad (2.28)$$

with  $\nabla_{\bar{i}\bar{j}}^{(2)} := \nabla_{(\bar{i}}^{(1)\gamma} \nabla_{\gamma\bar{j})}^{(1)}$  and  $\nabla_{\alpha\beta}^{(2)} := \nabla_{(\alpha}^{(1)\bar{k}} \nabla_{\beta)\bar{k}}^{(1)}$ . The action (2.26) is independent of the representative  $C_L^{(-4)}$  in the sense that it does not change under an arbitrary

---

<sup>3</sup>The operator  $\Delta_L^{(4)}$  is a covariant projection operator. Given a covariant left projective multiplet  $Q_L^{(n)}(v_L)$  of weight  $n$ , it may be represented in the form  $Q_L^{(n)} = \Delta_L^{(4)} T_L^{(n-4)}$ , for some left isotwistor superfield  $T_L^{(n-4)}(v_L)$ , see [14] for details.



infinitesimal variation of  $C_L^{(-4)}$  subject to the above conditions. The structure of  $S_R$  is analogous.

There are two equivalent action functionals to describe the dynamics of a single self-dual Abelian vector multiplet coupled to conformal supergravity. One of them is a right action formulated in terms of a right  $\mathcal{O}(2)$  multiplet  $G_R^{(2)}(v_R) = v_{\bar{i}}v_{\bar{j}}G^{\bar{i}\bar{j}}$ , which is associated with the superfield strength  $G^{\bar{i}\bar{j}}$  of the vector multiplet. This action, has the form<sup>4</sup> [14]

$$S_{\text{VM}}^{(+)} := \frac{\sqrt{2}}{\pi} \oint (v_R, dv_R) \int d^3|8 z E C_R^{(-4)} G_R^{(2)} \ln \frac{G_R^{(2)}}{i\Upsilon_R^{(1)}\check{\Upsilon}_R^{(1)}} , \quad (2.29)$$

where the weight-one arctic multiplet  $\Upsilon_R^{(1)}$  and its smile conjugate  $\check{\Upsilon}_R^{(1)}$  are pure gauge degrees of freedom. The action (2.29) is the 3D  $\mathcal{N} = 4$  counterpart of the projective-superspace action [29] for the 4D  $\mathcal{N} = 2$  improved tensor multiplet [30]. The other representation for  $S_{\text{VM}}^{(+)}$  makes use of a left tropical prepotential  $\mathcal{V}_L(v_L)$  of the self-dual vector multiplet with gauge transformations

$$\delta\mathcal{V}_L = \lambda_L + \check{\lambda}_L . \quad (2.30)$$

The gauge parameter  $\lambda_L$  is an arbitrary left arctic multiplet of weight zero. The gauge invariant field strength,  $G^{\bar{i}\bar{j}}$ , is related to  $\mathcal{V}_L$  through

$$G_R^{(2)}(v_R) = v_{\bar{i}}v_{\bar{j}}G^{\bar{i}\bar{j}} = \frac{i}{4}v_{\bar{i}}v_{\bar{j}} \oint \frac{(v_L, dv_L)}{2\pi} \frac{u_i u_j}{(v_L, u_L)^2} \nabla^{\alpha\bar{i}\bar{i}} \nabla_{\alpha}{}^{j\bar{j}} \mathcal{V}_L(v_L) . \quad (2.31)$$

Here  $u_L = u^i$  is a constant isospinor such that  $(v_L, u_L) \neq 0$  along the closed integration contour.<sup>5</sup> The action (2.29) can be recast as a left  $BF$ -type action [12]

$$S_{\text{VM}}^{(+)} = -\frac{1}{2\pi} \oint (v_L, dv_L) \int d^3|8 z E C_L^{(-4)} \mathcal{V}_L \mathbf{G}_L^{(2)} , \quad (2.32)$$

where  $\mathbf{G}_L^{(2)}(v_L) = v_i v_j \mathbf{G}^{ij}$  is the composite left  $\mathcal{O}(2)$  multiplet [12]

$$\begin{aligned} \mathbf{G}_L^{(2)} &= -\frac{i}{\sqrt{2}} v_i v_j \oint \frac{(v_R, dv_R)}{2\pi} \frac{u_{\bar{i}} u_{\bar{j}}}{(v_R, u_R)^2} \nabla^{\alpha\bar{i}\bar{i}} \nabla_{\alpha}{}^{j\bar{j}} \ln \frac{G_R^{(2)}}{i\Upsilon_R^{(1)}\check{\Upsilon}_R^{(1)}} \\ &= \frac{i}{4} v_i v_j \nabla^{\alpha\bar{i}\bar{i}} \nabla_{\alpha}{}^{j\bar{j}} \left( \frac{G_{\bar{i}\bar{j}}}{G_+} \right) . \end{aligned} \quad (2.33)$$

The composite left superfield  $\mathbf{G}^{ij}$  can be equivalently realised as the anti-self-dual  $\text{SO}(4)$  bivector  $\mathbf{G}_{-}^{IJ}$ .

---

<sup>4</sup>We should emphasise that in this paper we have defined the vector multiplet actions with “wrong” sign, because in our approach they correspond to superconformal compensators.

<sup>5</sup>One may show that the right-hand side of (2.31) is independent of  $u_L$ .

Similarly, the action for the anti-self-dual vector multiplet [14] can be recast as the right  $BF$ -type action [12]

$$S_{\text{VM}}^{(-)} := -\frac{1}{2\pi} \oint (v_{\text{R}}, dv_{\text{R}}) \int d^{3|8}z E C_{\text{R}}^{(-4)} \mathcal{V}_{\text{R}} \mathbf{G}_{\text{R}}^{(2)} , \quad (2.34)$$

where  $\mathbf{G}_{\text{R}}^{(2)}(v_{\text{R}}) = v_{\bar{i}} v_{\bar{j}} \mathbf{G}^{\bar{i}\bar{j}}$  is the composite right  $\mathcal{O}(2)$  multiplet [12]

$$\begin{aligned} \mathbf{G}_{\text{R}}^{(2)} &= -\frac{\mathbf{i}}{\sqrt{2}} v_{\bar{i}} v_{\bar{j}} \oint \frac{(v_{\text{L}}, dv_{\text{L}})}{2\pi} \frac{u_i u_j}{(v_{\text{L}}, u_{\text{L}})^2} \nabla^{\alpha\bar{i}\bar{i}} \nabla_{\alpha}{}^{j\bar{j}} \ln \frac{G_{\text{L}}^{(2)}}{\mathbf{i}\Upsilon_{\text{L}}^{(1)} \bar{\Upsilon}_{\text{L}}^{(1)}} \\ &= v_{\bar{i}} v_{\bar{j}} \frac{\mathbf{i}}{4} \nabla^{\alpha\bar{i}\bar{i}} \nabla_{\alpha}{}^{j\bar{j}} \left( \frac{G_{ij}}{G_{-}} \right) , \end{aligned} \quad (2.35)$$

and  $\mathcal{V}_{\text{R}}(v_{\text{R}})$  is the tropical prepotential of the anti-self-dual vector multiplet. The composite right superfield (2.35) can be equivalently realised as the self-dual  $\text{SO}(4)$  bivector  $\mathbf{G}_{+}^{IJ}$ .

The composite  $\mathcal{O}(2)$  multiplets can be expressed in terms of  $\text{SO}(4)$  vector indices as follows [12]

$$\mathbf{G}_{\pm}^{IJ} = X_{\mp}^{IJ} \pm \frac{1}{2} \varepsilon^{IJKL} X_{\mp KL} , \quad \frac{1}{2} \varepsilon_{IJKL} \mathbf{G}_{\pm}^{KL} = \pm \mathbf{G}_{\pm IJ} , \quad (2.36)$$

where we have defined

$$X_{\pm}^{IJ} := \frac{\mathbf{i}}{6G_{\pm}} \nabla^{\gamma[I} \nabla_{\gamma K} G_{\pm}^{J]K} + \frac{2\mathbf{i}}{9G_{\pm}^3} \nabla^{\alpha P} G_{\pm KP} \nabla_{\alpha Q} G_{\pm}^{Q[I} G_{\pm}^{J]K} . \quad (2.37)$$

To show that  $\mathbf{G}_{\pm}^{IJ}$  is primary and satisfies the Bianchi identity, the following identities prove useful

$$G_{\pm}^{IK} G_{\pm JK} = \frac{1}{2} \delta_J^I G_{\pm}^2 , \quad (2.38a)$$

$$\varepsilon^{IJKL} G_{\pm LP} = \mp 3 \delta_P^{[I} G_{\pm}^{JK]} . \quad (2.38b)$$

It is worth mentioning that the two  $\mathcal{N} = 4$  linear multiplet actions (2.32) and (2.34) are universal [12] in the sense that all known off-shell supergravity-matter systems (with the exception of pure conformal supergravity) may be described using such actions with appropriately engineered composite  $\mathcal{O}(2)$  multiplets  $\mathbf{G}_{\text{L}}^{(2)}$  and  $\mathbf{G}_{\text{R}}^{(2)}$ .

### 3 Minimal topologically massive supergravity

In this section we present two new supergravity-matter systems as models for minimal topologically massive supergravity.

### 3.1 Action principle and equations of motion

Our models for minimal topologically massive supergravity are described by  $\mathcal{N} = 4$  conformal supergravity coupled to a vector multiplet, either self-dual or anti-self-dual, via the following supergravity-matter actions:

$$\kappa S_{\pm} := \frac{1}{\mu} S_{\text{CSG}} + S_{\text{VM}}^{(\pm)} , \quad (3.1)$$

where  $S_{\text{CSG}}$  denotes the conformal supergravity action [17]. We will refer to the theories with actions  $S_+$  and  $S_-$  as the self-dual and anti-self-dual topologically massive supergravity (TMSG) theories, respectively.

As shown in [12], the equation of motion for the vector multiplet in the action (3.1) is equivalent to

$$\mathbf{G}_{\mp}{}^{IJ} = 0 , \quad (3.2)$$

while the equation of motion for the conformal supergravity multiplet (that is, the  $\mathcal{N} = 4$  Weyl supermultiplet) is

$$\frac{1}{\mu} W + T_{\pm} = 0 , \quad (3.3)$$

where  $T_{\pm}$  is the supercurrent given by

$$T_{\pm} = \pm G_{\pm} . \quad (3.4)$$

One can check that the supercurrent  $T_{\pm}$  obeys the conservation equation [31]

$$\nabla^{\alpha(I} \nabla_{\alpha}^{J)} T_{\pm} = \frac{1}{4} \delta^{IJ} \nabla_K^{\alpha} \nabla_{\alpha}^K T \quad (3.5)$$

when the matter equation of motion (3.2) is satisfied.

Making use of the Bianchi identity (2.7) as well as the equations of motion (3.2)–(3.4), one finds the following equations on  $G_{\pm}$ :

$$\left( \nabla^{\gamma(I} \nabla_{\gamma}^{J)} - \frac{1}{4} \delta^{IJ} \nabla_K^{\gamma} \nabla_{\gamma}^K \right) G_{\pm} = 0 , \quad (3.6a)$$

$$\left( \nabla_K^{\gamma} \nabla_{\gamma}^K \mp 8iW \right) G_{\pm}^{-1} = 0 , \quad (3.6b)$$

$$\frac{1}{\mu} W \pm G_{\pm} = 0 , \quad (3.6c)$$

$$\nabla_{(\alpha}^{[I} \nabla_{\beta}^{J]} G_{\pm}^{-1} = \pm \frac{1}{2} \varepsilon^{IJKL} \nabla_{(\alpha K} \nabla_{\beta) L} G_{\pm}^{-1} . \quad (3.6d)$$

We now turn to an analysis of the consequences of the equations of motion (3.6).

### 3.2 Analysing the equations of motion

To analyse the equations of motion corresponding to the action (3.1) we need to fix the gauge. Firstly, we use the special conformal transformations to make the dilatation connection vanish,  $B_A = 0$ . This corresponds to degauging of conformal superspace to  $\text{SO}(4)$  superspace [14] and gives rise to new torsion terms<sup>6</sup> which can be expressed in terms of superfields  $\mathcal{S}^{IJ}$ ,  $\mathcal{S}$ ,  $C_a^{IJ}$  and their covariant derivatives. We refer the reader to [14] for details and provide a summary of the salient details of  $\text{SO}(4)$  superspace in Appendix B.

Upon imposing the gauge  $B_A = 0$  one can show that (3.6) is equivalent to

$$\left(\mathcal{D}^{\gamma(I}\mathcal{D}^{J)}_{\gamma} - \frac{1}{4}\delta^{IJ}\mathcal{D}^{\gamma}_K\mathcal{D}^K_{\gamma} - 4i\mathcal{S}^{IJ}\right)G_{\pm} = 0, \quad (3.7a)$$

$$\left(\mathcal{D}^{\gamma}_K\mathcal{D}^K_{\gamma} + 8i(2\mathcal{S} \mp W)\right)G_{\pm}^{-1} = 0, \quad (3.7b)$$

$$\frac{1}{\mu}W \pm G_{\pm} = 0, \quad (3.7c)$$

$$(\mathcal{D}^{[I}_{(\alpha}\mathcal{D}^{J]}_{\beta)} - 4iC_{\alpha\beta}^{IJ})G_{\pm}^{-1} = \pm\frac{1}{2}\varepsilon^{IJKL}(\mathcal{D}_{(\alpha K}\mathcal{D}_{\beta)L} - 4iC_{\alpha\beta KL})G_{\pm}^{-1}, \quad (3.7d)$$

where  $\mathcal{D}^I_{\alpha}$  is the  $\text{SO}(4)$  superspace covariant derivative [14, 27] (see also [18]). In isospinor index notation, for the self-dual vector multiplet one obtains

$$\left(\mathcal{D}^{\gamma\bar{i}\bar{i}}\mathcal{D}_{\gamma\bar{i}\bar{i}} + 8i(2\mathcal{S} - W)\right)G_+^{-1} = 0, \quad (3.8a)$$

$$(\mathcal{D}^{(\bar{i}\bar{k}}_{\alpha}\mathcal{D}^{j)}_{\beta\bar{k}} - 4iC_{\alpha\beta}^{ij})G_+^{-1} = 0, \quad (3.8b)$$

$$(\mathcal{D}^{\gamma(i\bar{i}}\mathcal{D}^{j)\bar{j}}_{\gamma} - 4i\mathcal{S}^{ij\bar{j}\bar{j}})G_+ = 0, \quad (3.8c)$$

$$W + \mu G_+ = 0, \quad (3.8d)$$

while for the anti-self-dual vector multiplet one finds

$$\left(\mathcal{D}^{\gamma\bar{i}\bar{i}}\mathcal{D}_{\gamma\bar{i}\bar{i}} + 8i(2\mathcal{S} + W)\right)G_-^{-1} = 0, \quad (3.9a)$$

$$(\mathcal{D}^{k(\bar{i}}_{\alpha}\mathcal{D}^{j)}_{\beta k} - 4iC_{\alpha\beta}^{ij})G_-^{-1} = 0, \quad (3.9b)$$

$$(\mathcal{D}^{\gamma(i\bar{i}}\mathcal{D}^{j)\bar{j}}_{\gamma} - 4i\mathcal{S}^{ij\bar{j}\bar{j}})G_- = 0, \quad (3.9c)$$

$$W - \mu G_- = 0. \quad (3.9d)$$

One should keep in mind that the equations of motion for  $G_+$  and  $G_-$  derived from the actions  $S_+$  and  $S_-$ , respectively, were used in the above results.

---

<sup>6</sup>See [17] for more details. It is important to note that the  $\text{SO}(4)$  connection of  $\text{SO}(4)$  superspace differs from the one from the one of conformal superspace by a redefinition, for details see [18].

Under super-Weyl transformations the  $SO(4)$ -covariant derivatives and the torsion terms transform as<sup>7</sup>

$$\mathcal{D}_\alpha^I \rightarrow \mathcal{D}'_\alpha^I = e^{\frac{1}{2}\sigma} \left( \mathcal{D}_\alpha^I + (\mathcal{D}^{\beta I} \sigma) M_{\alpha\beta} + (\mathcal{D}_{\alpha J} \sigma) N^{IJ} \right), \quad (3.10a)$$

$$\mathcal{S}^{IJ} \rightarrow \mathcal{S}'^{IJ} = \frac{i}{4} e^{2\sigma} (\mathcal{D}^{\gamma(I} \mathcal{D}_\gamma^{J)} - \frac{1}{4} \delta^{IJ} \mathcal{D}^{\gamma K} \mathcal{D}_{\gamma K} - 4i \mathcal{S}^{IJ}) e^{-\sigma}, \quad (3.10b)$$

$$\mathcal{S} \rightarrow \mathcal{S}' = -\frac{i}{16} (\mathcal{D}_K^\gamma \mathcal{D}_\gamma^K + 16i \mathcal{S}) e^\sigma, \quad (3.10c)$$

$$C_a'^{IJ} \rightarrow C_a^{IJ} = -\frac{i}{8} (\gamma_a)^{\alpha\beta} (\mathcal{D}_\alpha^{[I} \mathcal{D}_\beta^{J]} - 4i C_{\alpha\beta}^{IJ}) e^\sigma, \quad (3.10d)$$

$$W \rightarrow W' = e^\sigma W, \quad (3.10e)$$

where  $\sigma$  is a real unconstrained superfield. Within the superconformal framework, all supergravity-matter actions are required to be super-Weyl invariant.

The super-Weyl gauge freedom may be used to impose useful gauge conditions. For instance, one can make use of the super-Weyl transformations to gauge away the self-dual or anti-self-dual part of  $C_a^{IJ}$  such that the remaining torsion components are expressed directly in terms of the matter fields. For concreteness, let us consider the theory described by the action  $S_+$ , with corresponding equations of motion (3.8), and gauge away  $C_a^{\bar{i}\bar{j}}$  via a super-Weyl transformation. We then find

$$W = -\mu G_+, \quad (3.11a)$$

$$\mathcal{S}^{ij\bar{i}\bar{j}} = -\frac{i}{4} G_+^{-1} \mathcal{D}^{\gamma(i\bar{i}} \mathcal{D}_\gamma^{j)\bar{j}} G_+, \quad (3.11b)$$

$$2\mathcal{S} - W = \frac{i}{8} G_+ \mathcal{D}^{\gamma\bar{i}\bar{i}} \mathcal{D}_{\gamma\bar{i}\bar{i}} G_+^{-1}, \quad (3.11c)$$

$$C_{\alpha\beta}^{ij} = -\frac{i}{4} G_+ \mathcal{D}_\alpha^{(i\bar{k}} \mathcal{D}_{\beta}^{j)\bar{k}} G_+^{-1}, \quad (3.11d)$$

$$C_a^{\bar{i}\bar{j}} = 0. \quad (3.11e)$$

In this gauge, we see that the geometry is determined in terms of a single superfield, which is chosen to be the scalar  $G_+$ . After imposing this super-Weyl gauge condition it is possible to show that there is enough super-Weyl freedom left to impose the additional condition

$$2\mathcal{S} + W = 0, \quad (3.12)$$

see Appendix C for the derivation. This condition proves to lead to the following nonlinear equation for  $G_+$ :

$$\mathcal{D}^{\gamma\bar{i}\bar{i}} \mathcal{D}_{\gamma\bar{i}\bar{i}} G_+^{-1} + 16i\mu = 0. \quad (3.13)$$

---

<sup>7</sup>The infinitesimal form was given in [14, 25].

The main virtue of the super-Weyl gauge conditions imposed is that all the torsion and curvature tensors are descendants of the single scalar superfield  $G_+$ . However, this gauge choice is not particularly useful from the point of view of studying (maximally) supersymmetric backgrounds. A more convenient super-Weyl gauge fixing is  $G_+ = \text{const}$ . We spell out the implications of such a gauge condition below.

Given a vector multiplet with a superfield strength  $G^{IJ}$  such that  $G$  is nowhere vanishing, one can always make use of the super-Weyl transformations to choose a gauge where

$$G = \frac{1}{2}G^{IJ}G_{IJ} = 1, \quad \mathcal{D}_\alpha^I G^{JK} = 0. \quad (3.14)$$

Such a gauge condition has slightly different consequences on the superspace geometry for the two vector multiplets  $G_+^{IJ}$  and  $G_-^{IJ}$  satisfying the equations of motion (3.2) and (3.3). In both cases the super-Cotton tensor is constant,

$$W = \text{const} \implies \mathcal{S}^{IJ} = 0, \quad (3.15)$$

while the constraints on the remaining torsion components differ. For the on-shell self-dual vector multiplet one finds the following consistency conditions

$$\frac{1}{2}\varepsilon_{IJKL}C_a^{KL} = C_{aIJ}, \quad 2\mathcal{S} - W = 0, \quad (3.16)$$

while for the on-shell anti-self-dual vector multiplet one finds

$$-\frac{1}{2}\varepsilon_{IJKL}C_a^{KL} = C_{aIJ}, \quad 2\mathcal{S} + W = 0. \quad (3.17)$$

In the case where  $C_a^{IJ}$  vanishes, the algebra of covariant derivatives coincides with that of  $(4, 0)$  AdS superspace in the critical case where  $2\mathcal{S} \mp W = 0$ , see [25].<sup>8</sup> In general, however,  $C_a^{IJ}$  does not vanish and instead satisfies some differential conditions implied by the Bianchi identities

$$\begin{aligned} & [[\mathcal{D}_A, \mathcal{D}_B], \mathcal{D}_C] + (-1)^{\varepsilon_A(\varepsilon_B + \varepsilon_C)} [[\mathcal{D}_B, \mathcal{D}_C], \mathcal{D}_A] \\ & + (-1)^{\varepsilon_C(\varepsilon_A + \varepsilon_B)} [[\mathcal{D}_C, \mathcal{D}_A], \mathcal{D}_B] = 0. \end{aligned} \quad (3.18)$$

To analyse the Bianchi identities in detail it will be useful to convert to isospinor notation.

We consider in detail the self-dual TMSG theory. In the isospinor notation, the covariant derivative algebra which follows from the equations of motion is

$$\{\mathcal{D}_\alpha^{\bar{i}\bar{j}}, \mathcal{D}_\beta^{\bar{j}\bar{j}}\} = 2i\varepsilon^{ij}\varepsilon^{\bar{i}\bar{j}}\mathcal{D}_{\alpha\beta} + 4i\varepsilon_{\alpha\beta}\varepsilon^{\bar{i}\bar{j}}WL^{ij} + 4iC_{\alpha\beta}^{\bar{i}\bar{j}}L^{ij}$$

---

<sup>8</sup>The  $\mathcal{N} = 4$  super-Cotton tensor is denoted by  $X$  in [14, 25].

$$+2i\varepsilon_{\alpha\beta}\varepsilon^{ij}C^{\gamma\delta\bar{i}\bar{j}}M_{\gamma\delta}-2i\varepsilon^{ij}\varepsilon^{\bar{i}\bar{j}}WM_{\alpha\beta} . \quad (3.19a)$$

Analysing the Bianchi identities (3.18) determines the remainder of the covariant derivative algebra:

$$\begin{aligned} [\mathcal{D}_{\alpha\beta}, \mathcal{D}_{\gamma}^{k\bar{k}}] &= -\varepsilon_{\gamma(\alpha}W\mathcal{D}_{\beta)}^{k\bar{k}} + (\varepsilon_{\gamma(\alpha}C_{\beta)\delta}^{\bar{k}\bar{j}} + \varepsilon_{\delta(\alpha}C_{\beta)\gamma}^{\bar{k}\bar{j}})\mathcal{D}^{\delta k}_{\bar{j}} \\ &\quad + 2\varepsilon_{\gamma(\alpha}C_{\beta)\delta\rho}^{k\bar{k}}M^{\delta\rho} - 2C_{\alpha\beta\gamma}^{j\bar{k}}L_j^k , \end{aligned} \quad (3.19b)$$

$$\begin{aligned} [\mathcal{D}_{\alpha\beta}, \mathcal{D}_{\gamma\delta}] &= i\varepsilon_{\gamma(\alpha}C_{\beta)\delta\rho k\bar{k}}\mathcal{D}^{\rho k\bar{k}} + i\varepsilon_{\delta(\alpha}C_{\beta)\gamma\rho k\bar{k}}\mathcal{D}^{\rho k\bar{k}} \\ &\quad + \varepsilon_{\delta(\alpha}W^2M_{\beta)\gamma} + \varepsilon_{\gamma(\alpha}W^2M_{\beta)\delta} \\ &\quad + \frac{i}{12}\varepsilon_{\delta(\alpha}(\mathcal{D}_{\beta)}^{k\bar{k}}\mathcal{D}_{\gamma k}^{\bar{l}}C_{\rho\sigma\bar{k}\bar{l}})M^{\rho\sigma} + \frac{i}{12}\varepsilon_{\gamma(\alpha}(\mathcal{D}_{\beta)}^{k\bar{k}}\mathcal{D}_{\delta k}^{\bar{l}}C_{\rho\sigma\bar{k}\bar{l}})M^{\rho\sigma} \\ &\quad - \varepsilon_{\delta(\alpha}C_{\beta)\gamma\bar{k}\bar{l}}C^{\rho\sigma k\bar{l}}M_{\rho\sigma} - \varepsilon_{\gamma(\alpha}C_{\beta)\delta\bar{k}\bar{l}}C^{\rho\sigma k\bar{l}}M_{\rho\sigma} , \end{aligned} \quad (3.19c)$$

as well as the following differential constraint on  $C_a^{\bar{i}\bar{j}}$

$$\mathcal{D}_{\alpha}^{\bar{i}\bar{i}}C_{\beta\gamma}^{\bar{j}\bar{k}} = 2\varepsilon^{\bar{i}\bar{j}}C_{\alpha\beta\gamma}^{i\bar{k}} . \quad (3.20)$$

The above constraint implies, in turn,

$$\mathcal{D}_{\alpha}^{\gamma}C_{\beta\gamma}^{\bar{i}\bar{j}} + C_{(\alpha}^{\gamma}{}_{\bar{k}}(\bar{i}C_{\beta)\gamma}^{\bar{j}\bar{k}}) + 2WC_{\alpha\beta}^{\bar{i}\bar{j}} = 0 . \quad (3.21)$$

Since the  $SU(2)_R$  curvature vanishes, we can completely gauge away the corresponding connection. Such a gauge condition is assumed in what follows. In this gauge, the field strength  $G^{\bar{i}\bar{j}}$  becomes a constant symmetric isospinor subject to the normalisation condition  $G^{\bar{i}\bar{j}}G_{\bar{i}\bar{j}} = 1$ . It is invariant under a  $U(1)$  subgroup of  $SU(2)_R$ .

We are now in a position to describe all maximally supersymmetric solutions of the theory. In accordance with the general superspace analysis of supersymmetric backgrounds in diverse dimensions [32, 33, 34], such superspaces have to comply with the additional constraint

$$\mathcal{D}_{\alpha}^{\bar{i}\bar{i}}C_{\beta\gamma}^{\bar{j}\bar{k}} = 0 , \quad (3.22)$$

which leads to the integrability conditions

$$(\mathcal{D}_a - WM_a)C_b^{\bar{j}\bar{k}} = 0 , \quad (3.23a)$$

$$C^{\gamma}{}_{(\alpha}{}^{\bar{i}\bar{j}}C_{\beta)\gamma}^{\bar{k}\bar{l}} = 0 . \quad (3.23b)$$

The general solution of (3.23b) is

$$C_{\alpha\beta}^{\bar{i}\bar{j}} = C_{\alpha\beta}C^{\bar{i}\bar{j}} , \quad (3.24)$$

where  $C^{\bar{i}\bar{j}}$  is a constant symmetric rank-2 isospinor. Without loss of generality,  $C^{\bar{i}\bar{j}}$  can be normalised as  $C^{\bar{i}\bar{j}}C_{\bar{i}\bar{j}} = 1$ . The covariant constancy conditions (3.22) and (3.23a) now amount to

$$\mathcal{D}_\alpha^{\bar{i}\bar{i}}C_b = 0, \quad (\mathcal{D}_a - WM_a)C_b = 0. \quad (3.25)$$

We recall that the Lorentz generator with a vector index,  $M_a$ , acts on a three-vector by the rule  $M_a C_b = \varepsilon_{abc} C^c$ . The second condition in (3.25) implies that  $C_b$  is a Killing vector of constant norm,

$$\mathcal{D}_a C_b + \mathcal{D}_b C_a = 0, \quad C^2 = C^a C_a = \text{const}. \quad (3.26)$$

Thus there are three types of backgrounds depending on whether the Killing vector  $C^a$  is chosen to be time-like, space-like or null. The algebra of covariant derivatives for such a background is

$$\begin{aligned} \{\mathcal{D}_\alpha^{\bar{i}\bar{i}}, \mathcal{D}_\beta^{\bar{j}\bar{j}}\} &= 2i\varepsilon^{ij}\varepsilon^{\bar{i}\bar{j}}(\mathcal{D}_{\alpha\beta} - WM_{\alpha\beta}) + 4i\varepsilon_{\alpha\beta}\varepsilon^{\bar{i}\bar{j}}WL^{ij} + 4iC^{\bar{i}\bar{j}}C_{\alpha\beta}L^{ij} \\ &\quad + 2i\varepsilon_{\alpha\beta}\varepsilon^{ij}C^{\bar{i}\bar{j}}C^{\gamma\delta}M_{\gamma\delta}, \end{aligned} \quad (3.27a)$$

$$[\mathcal{D}_{\alpha\beta}, \mathcal{D}_\gamma^{k\bar{k}}] = -\varepsilon_{\gamma(\alpha}W\mathcal{D}_{\beta)}^{k\bar{k}} + (\varepsilon_{\gamma(\alpha}C_{\beta)\delta}^{\bar{k}\bar{j}} + \varepsilon_{\delta(\alpha}C_{\beta)\gamma}^{\bar{k}\bar{j}})\mathcal{D}^{\delta k}_{\bar{j}}, \quad (3.27b)$$

$$[\mathcal{D}_{\alpha\beta}, \mathcal{D}_{\gamma\delta}] = W^2(\varepsilon_{\delta(\alpha}M_{\beta)\gamma} + \varepsilon_{\gamma(\alpha}M_{\beta)\delta}) - (\varepsilon_{\delta(\alpha}C_{\beta)\gamma} + \varepsilon_{\gamma(\alpha}C_{\beta)\delta})C^{\rho\sigma}M_{\rho\sigma}. \quad (3.27c)$$

One may think of this algebra as a Lie superalgebra.<sup>9</sup> By construction, the theory involves the constant symmetric isospinor  $G^{\bar{i}\bar{j}}$  being invariant under a U(1) subgroup of the group SU(2)<sub>R</sub>. If  $C^{\bar{i}\bar{j}}$  does not coincide with  $G^{\bar{i}\bar{j}}$ , then the group SU(2)<sub>R</sub> is completely broken. This indicates that  $C^{\bar{i}\bar{j}} = G^{\bar{i}\bar{j}}$ .

The simplest maximally supersymmetric solution of the theory is characterised by (see also [20])

$$C_a^{\bar{i}\bar{j}} = 0. \quad (3.28)$$

It corresponds to the critical (4,0) AdS superspace introduced in [25]. Its algebra of covariant derivatives is as follows:

$$\{\mathcal{D}_\alpha^{\bar{i}\bar{i}}, \mathcal{D}_\beta^{\bar{j}\bar{j}}\} = 2i\varepsilon^{ij}\varepsilon^{\bar{i}\bar{j}}(\mathcal{D}_{\alpha\beta} - WM_{\alpha\beta}) + 4i\varepsilon_{\alpha\beta}\varepsilon^{\bar{i}\bar{j}}WL^{ij}, \quad (3.29a)$$

$$[\mathcal{D}_a, \mathcal{D}_\beta^{\bar{j}\bar{j}}] = \frac{1}{2}W(\gamma_a)_\beta^\gamma \mathcal{D}_\gamma^{\bar{j}\bar{j}}, \quad [\mathcal{D}_a, \mathcal{D}_b] = -W^2 M_{ab}. \quad (3.29b)$$

The last relation shows that the cosmological constant is  $\Lambda = -W^2 = -\ell^{-2}$ , in agreement with [20, 25]. Here  $\ell$  is the radius of curvature in AdS<sub>3</sub>. The latter relation is equivalent to  $\mu\ell = 1$ , which corresponds to chiral gravity [24].

---

<sup>9</sup>More precisely, (3.27) is isomorphic to the Lie superalgebra corresponding to the isometry supergroup of the background superspace under consideration.



More generally, the  $(p, q)$  AdS superspaces,  $p + q = \mathcal{N}$ , in three dimensions were classified in [25]<sup>10</sup>. In the  $\mathcal{N} = 4$  case, the (3,1) and (2,2) AdS superspaces are necessarily conformally flat,  $W = 0$ . The distinguished feature of (4,0) AdS supersymmetry is that the super-Cotton scalar  $W$  may have a non-zero value. The algebra of covariant derivatives is given by [25]

$$\{\mathcal{D}_\alpha^{\bar{i}\bar{j}}, \mathcal{D}_\beta^{j\bar{j}}\} = 2i\varepsilon^{ij}\varepsilon^{\bar{i}\bar{j}}\mathcal{D}_{\alpha\beta} + 2i\varepsilon_{\alpha\beta}\varepsilon^{\bar{i}\bar{j}}(2\mathcal{S} + W)L^{ij} + 2i\varepsilon_{\alpha\beta}\varepsilon^{ij}(2\mathcal{S} - W)R^{\bar{i}\bar{j}} - 4i\mathcal{S}\varepsilon^{ij}\varepsilon^{\bar{i}\bar{j}}M_{\alpha\beta}, \quad (3.30a)$$

$$[\mathcal{D}_a, \mathcal{D}_\beta^{j\bar{j}}] = \mathcal{S}(\gamma_a)_\beta{}^\gamma \mathcal{D}_\gamma^{j\bar{j}}, \quad [\mathcal{D}_a, \mathcal{D}_b] = -4\mathcal{S}^2 M_{ab}, \quad (3.30b)$$

where the positive constant  $\mathcal{S}$  determines the curvature of  $\text{AdS}_3$ . For a generic value of  $W$  the entire  $\text{SO}(4)$   $R$ -symmetry group belongs to the superspace holonomy group. But there are two points in which either the  $\text{SU}(2)_R$  or the  $\text{SU}(2)_L$  curvature vanishes and the structure group is reduced. These are given by

$$W = \pm 2\mathcal{S} \quad (3.31)$$

and correspond to the critical (4,0) AdS superspaces. As briefly discussed in [35], the isometry group of (4,0) AdS superspace is isomorphic to  $\text{D}(2, 1; \alpha) \times \text{SL}(2, \mathbb{R})$  in the non-critical case  $W \neq \pm 2\mathcal{S}$ , where  $\text{D}(2, 1; \alpha)$  is one of the exceptional simple supergroups, with the real number  $\alpha \neq -1, 0$ , see e.g. [36, 37] for reviews. The supergroup parameter  $\alpha$  is related to the (4,0) AdS parameter  $q = 1 + \frac{W}{2\mathcal{S}}$  introduced in [35]. If the values of  $\alpha$  are restricted to the range<sup>11</sup>  $-1 < \alpha \leq -\frac{1}{2}$ , then we can identify  $-2\alpha = 1 + \frac{W}{2\mathcal{S}}$ . The case  $\alpha = -\frac{1}{2}$  corresponds to the conformally flat (4,0) AdS superspace, for which  $W = 0$ . Its isometry group is  $\text{OSp}(4|2) \times \text{SL}(2, \mathbb{R})$ . The limiting choice  $\alpha = -1$  corresponds to one of the two critical (4,0) AdS cases,  $W = 2\mathcal{S}$ .<sup>12</sup> The isometry group of this (4,0) AdS superspace is  $\text{SU}(1, 1|2) \rtimes \text{SU}(2) \times \text{SL}(2, \mathbb{R})$ , see also the discussion in [38].

If  $C^a \neq 0$ , the maximally supersymmetric background (3.27) describes a warped critical (4,0) AdS superspace. The bosonic body of such a superspace is warped  $\text{AdS}_3$  spacetime associated with the Killing vector  $c^a(x) = C^a(z)|_{\theta=0}$ . Warped  $\text{AdS}_3$  spacetimes have been discussed in detail in the literature, see [39, 40, 41, 42] and references therein. In the  $\mathcal{N} = 2$  supersymmetric case, the (super)space geometry of maximally supersymmetric warped (1,1) and (2,0) AdS backgrounds was described in

<sup>10</sup>In three dimensions,  $\mathcal{N}$ -extended AdS supergravity exists in several incarnations [7] known as the  $(p, q)$  AdS supergravity theories, where the integers  $p \geq q \geq 0$  are such that  $\mathcal{N} = p + q$ .

<sup>11</sup>Not all values of  $\alpha$  lead to distinct supergroups, since the supergroups defined by the parameters  $\alpha^{\pm 1}$ ,  $-(1 + \alpha)^{\pm 1}$  and  $-\alpha^{\pm 1}(1 + \alpha)^{\mp 1}$  are isomorphic [36, 37].

<sup>12</sup>The isometry groups of the two critical (4,0) AdS superspaces are isomorphic.

[11] and further elaborated in [34]. Supersymmetric warped (1,1) AdS backgrounds were thoroughly studied in [43].

We now linearise the equation (3.21) around the critical (4,0) AdS superspace and let  $C_a^{\bar{i}\bar{j}} = \delta C_a^{\bar{i}\bar{j}}$  where  $\delta C_a^{\bar{i}\bar{j}}$  is a small disturbance. Eq. (3.21) turns into

$$\mathcal{D}_\alpha{}^\gamma \delta C_{\beta\gamma}^{\bar{i}\bar{j}} - 2\mu \delta C_{\alpha\beta}^{\bar{i}\bar{j}} = 0 \implies \mathcal{D}^a \delta C_a^{\bar{i}\bar{j}} = 0, \quad (3.32)$$

where  $\mathcal{D}_a$  denotes the vector covariant derivative of the critical (4,0) AdS superspace. After applying another vector derivative one finds the equation

$$(\mathcal{D}^a \mathcal{D}_a - 2\mu^2) \delta C_b^{\bar{i}\bar{j}} = 0. \quad (3.33)$$

One can also derive further equations on descendants of  $\delta C_{\alpha\beta}^{\bar{i}\bar{j}}$  using the constraint (3.20). In particular, one finds

$$(\mathcal{D}_\alpha{}^\delta - \frac{3}{2}\mu\delta_\alpha^\delta) \delta C_{\beta\gamma\delta}^{\bar{i}\bar{i}} = 0, \quad \delta C_{\alpha\beta\gamma}^{\bar{i}\bar{i}} := \frac{1}{3} \mathcal{D}_\alpha{}^{\bar{i}}{}_{\bar{j}} \delta C_{\beta\gamma}^{\bar{i}\bar{j}}, \quad (3.34a)$$

$$(\mathcal{D}_\alpha{}^\rho - \mu\delta_\alpha^\rho) \delta C_{\beta\gamma\delta\rho} = 0, \quad \delta C_{\alpha\beta\gamma\delta} := \mathcal{D}_{(\alpha}^{\bar{i}\bar{i}} \delta C_{\beta\gamma\delta)}^{\bar{i}\bar{i}}, \quad (3.34b)$$

where  $\mathcal{D}_\alpha^{\bar{i}\bar{i}}$  denotes the spinor covariant derivative of the critical (4,0) AdS superspace. The component projection of  $\delta C_{\alpha\beta\gamma}^{\bar{i}\bar{i}}$  is proportional to the linearised gravitino field strength, while  $\delta C_{\alpha\beta\gamma\delta}$  is proportional to the linearised Cotton tensor. These superfields can be shown to satisfy the following consequences of eqs. (3.34):

$$(\mathcal{D}^a \mathcal{D}_a + \frac{1}{4}\mu^2) \delta C_{\alpha\beta\gamma}^{\bar{i}\bar{i}} = 0, \quad (3.35a)$$

$$(\mathcal{D}^a \mathcal{D}_a + 2\mu^2) \delta C_{\alpha\beta\gamma\delta} = 0. \quad (3.35b)$$

In the above we made use of the following result for a symmetric rank-(2s) superfield  $T_{\alpha_1 \dots \alpha_{2s}} = T_{(\alpha_1 \dots \alpha_{2s})}$  (with isospinor indices suppressed):

$$(\mathcal{D}_{\alpha_1}{}^\beta - \delta_{\alpha_1}{}^\beta \frac{\mu}{\eta}) T_{\alpha_2 \dots \alpha_{2s} \beta} = 0 \implies (\mathcal{D}^a \mathcal{D}_a - \frac{\mu^2}{\eta^2} + (s+1)\mu^2) T_{\alpha_1 \dots \alpha_{2s}} = 0, \quad (3.36)$$

with  $\eta$  a dimensionless parameter. Computing the bar-projection of the equations (3.32), (3.34a) and (3.34b), we can determine the representations of the AdS group  $SO(2,2)$  to which the fields  $\delta C_{\alpha\beta}^{\bar{i}\bar{j}}$ ,  $\delta C_{\alpha\beta\gamma}^{\bar{i}\bar{i}}$  and  $\delta C_{\alpha\beta\gamma\delta}$  belong. We recall that the unitary representations of  $SO(2,2)$ , denoted  $D(E_0, \hat{s})$ , are labelled by two real weights  $(E_0, \hat{s})$ , where  $E_0$  is the lowest energy and  $\hat{s}$  is the helicity, see e.g. [44]. The weights obey the unitarity bound  $E_0 \geq |\hat{s}|$  for  $\hat{s} > 0$ , where the representations with  $E_0 = |\hat{s}| > 0$  are called singleton representations. For a superfield  $T_{\alpha_1 \dots \alpha_{2s}}$  obeying the first-order equation (3.36), its lowest component  $T_{\alpha_1 \dots \alpha_{2s}}|$  transforms in the representation with

$$E_0 = 1 + \frac{1}{|\eta|} , \quad \hat{s} = \frac{s\eta}{|\eta|} , \quad (3.37)$$

as follows from the analysis in [44] (see also [45]). Thus the gravitational field  $\delta C_{\alpha\beta\gamma\delta}|$  is a helicity 2 singleton, while the spin-1 and spin-3/2 fields,  $\delta C_{\alpha\beta}{}^{\bar{i}\bar{j}}|$  and  $\delta C_{\alpha\beta\gamma}{}^{\bar{i}\bar{i}}|$ , are massive.

In the above we worked with the self-dual TMSG theory, however the analysis of the equations of motion corresponding to the action  $S_-$  is completely analogous. There one finds the covariant derivative algebra is

$$\begin{aligned} \{\mathcal{D}_{\alpha}{}^{\bar{i}\bar{i}}, \mathcal{D}_{\beta}{}^{j\bar{j}}\} &= 2i\varepsilon^{ij}\varepsilon^{\bar{i}\bar{j}}\mathcal{D}_{\alpha\beta} - 4i\varepsilon_{\alpha\beta}\varepsilon^{ij}WR^{\bar{i}\bar{j}} + 4iC_{\alpha\beta}{}^{ij}R^{\bar{i}\bar{j}} \\ &\quad + 2i\varepsilon_{\alpha\beta}\varepsilon^{\bar{i}\bar{j}}C^{\gamma\delta ij}M_{\gamma\delta} + 2i\varepsilon^{ij}\varepsilon^{\bar{i}\bar{j}}WM_{\alpha\beta} , \end{aligned} \quad (3.38)$$

where  $C_a{}^{ij}$  satisfies the Bianchi identity

$$\mathcal{D}_{\alpha}{}^{\bar{i}\bar{i}}C_{\beta\gamma}{}^{jk} = 2\varepsilon^{(j}C_{\alpha\beta\gamma}{}^{k)\bar{i}} . \quad (3.39)$$

Using the above equation one finds

$$\mathcal{D}_{\alpha}{}^{\gamma}C_{\beta\gamma}{}^{ij} + C_{(\alpha}{}^{\gamma}{}_k{}^{(i}C_{\beta)\gamma}{}^{j)k} - 2WC_{\alpha\beta}{}^{ij} = 0 . \quad (3.40)$$

The solution  $C_a{}^{ij} = 0$  corresponds to  $(4, 0)$  AdS superspace in the critical case  $2S = -W$ . We now linearise around the  $(4, 0)$  AdS superspace and set  $C_a{}^{ij} = \delta C_a{}^{ij}$  where  $\delta C_a{}^{ij}$  is a small disturbance. It can be seen that  $\delta C_a{}^{ij}$  obeys the equation

$$\mathcal{D}_{\alpha}{}^{\gamma}\delta C_{\beta\gamma}{}^{ij} - 2\mu\delta C_{\alpha\beta}{}^{ij} = 0 , \quad (3.41)$$

where  $\mathcal{D}_a$  corresponds to the vector covariant derivative of the  $(4, 0)$  AdS superspace. After applying another vector derivative one finds

$$(\mathcal{D}^a\mathcal{D}_a - 2\mu^2)\delta C_b{}^{ij} = 0 . \quad (3.42)$$

## 4 Component actions

In this section we give the component results corresponding to the minimal  $\mathcal{N} = 4$  topologically massive supergravity action (3.1).

## 4.1 The component conformal supergravity action

The complete component analysis of the  $\mathcal{N}$ -extended Weyl multiplet was given in [17]. Here we specialise to the  $\mathcal{N} = 4$  case where the auxiliary fields coming from the super-Cotton tensor are defined as:

$$w := \frac{1}{4!} \varepsilon_{IJKL} w^{IJKL} = W|, \quad y := \frac{1}{4!} \varepsilon_{IJKL} y^{IJKL} = -\frac{i}{4} \nabla_I^\alpha \nabla_\alpha^I W|, \quad (4.1a)$$

$$w_{\alpha L} := \frac{1}{3!} \varepsilon_{IJKL} w_\alpha^{IJK} = -\frac{i}{2} \nabla_{\alpha L} W|. \quad (4.1b)$$

The full  $\mathcal{N} = 4$  conformal supergravity action was given in [17] and is

$$\begin{aligned} S_{\text{CSG}} = \frac{1}{8} \int d^3x e \left\{ \varepsilon^{abc} (\omega_a^{fg} \mathcal{R}_{bcfg} - \frac{2}{3} \omega_{af}^g \omega_{bg}^h \omega_{ch}^f - \frac{i}{2} \Psi_{bcI}^\alpha (\gamma_d)_\alpha^\beta (\gamma_a)_\beta^\gamma \varepsilon^{def} \Psi_{ef\gamma}^I \right. \\ \left. - 2 \mathcal{R}_{ab}^{IJ} V_{cIJ} - \frac{4}{3} V_a^{IJ} V_{bI}^K V_{cKJ} \right) \\ \left. - 32i w_I^\alpha w_\alpha^I - 8wy - 16i \psi_a^\alpha (\gamma^a)_\alpha^\beta w_\beta^I w - 2i \varepsilon^{abc} (\gamma_a)_{\alpha\beta} \psi_b^\alpha \psi_c^\beta w^2 \right\}, \quad (4.2) \end{aligned}$$

where the component curvatures  $\mathcal{R}_{ab}^{cd}$  and  $\mathcal{R}_{ab}^{IJ}$  are defined as

$$\mathcal{R}_{ab}^{cd} := 2e_a^m e_b^n \partial_{[m} \omega_{n]}^{cd} - 2\omega_{[a}^{cf} \omega_{b]f}^d, \quad (4.3a)$$

$$\mathcal{R}_{ab}^{IJ} := 2e_a^m e_b^n \partial_{[m} V_{n]}^{IJ} - 2V_{[a}^{IK} V_{b]K}^J. \quad (4.3b)$$

## 4.2 The component vector multiplet actions

The component  $\mathcal{N} = 4$  linear multiplet actions were given in [12]. Making use of the results there, one can construct the left and right vector multiplet actions.

The component fields of the vector multiplets are defined as

$$g_\pm^{IJ} := G_\pm^{IJ}|, \quad (4.4a)$$

$$\lambda_{(\pm)\alpha}^I := \frac{2}{3} \nabla_{\alpha J} G_\pm^{IJ}|, \quad (4.4b)$$

$$h_{(\pm)}^{IJ} := \frac{i}{3} \nabla^{\gamma[I} \nabla_{\gamma K} G_\pm^{J]K}|, \quad (4.4c)$$

$$f_{(\pm)ab} := -\frac{i}{24} \varepsilon_{abc} (\gamma^c)^{\alpha\beta} \nabla_\alpha^K \nabla_\beta^L G_{\pm KL}| - \frac{1}{2} (\psi_{[a}^K \gamma_{b]} \lambda_{(\pm)K}) + \frac{i}{2} \psi_a^{\gamma K} \psi_{b\gamma}^L g_{\pm KL}, \quad (4.4d)$$

where  $g_\pm^{IJ}$  is (anti-)self-dual

$$\frac{1}{2} \varepsilon^{IJKL} g_{\pm KL} = \pm g_\pm^{IJ}. \quad (4.5)$$

The component gauge one-forms  $v_{(\pm)a}$  are defined as

$$v_{(\pm)a} := e_a^m v_{(\pm)m} , \quad f_{(\pm)ab} = 2e_a^m e_b^n \partial_{[m} v_{(\pm)n]} , \quad v_{(\pm)m} := V_{\pm m} | , \quad (4.6)$$

where  $V_{\pm}$  is the superspace gauge one-form associated with the field strength  $G_{\pm}^{IJ}$ .

It is useful to replace  $h_{(\pm)}^{IJ}$  by the fields

$$\begin{aligned} \hat{h}_{\pm}^{IJ} &= \frac{1}{2}(h_{(\mp)}^{IJ} + \tilde{h}_{(\mp)}^{IJ}) \\ &= h_{(\mp)}^{IJ} \mp 2wg_{\mp}^{IJ} , \end{aligned} \quad (4.7)$$

which proves to be (anti-)self-dual

$$\frac{1}{2}\varepsilon^{IJKL}\hat{h}_{\pm KL} = \pm \hat{h}_{\pm}^{IJ} . \quad (4.8)$$

The component self-dual vector multiplet action is

$$\begin{aligned} S_{\text{VM}}^{(+)} &= - \int d^3x e \left( \varepsilon^{abc} v_{(+a)} \mathbf{f}_{(+b)c} + \frac{1}{4} \hat{\mathbf{h}}_+^{IJ} g_{+IJ} + \frac{1}{4} \hat{\mathbf{h}}_-^{IJ} g_{-IJ} - \frac{i}{2} \lambda^{\alpha I} \lambda_{\alpha I} \right. \\ &\quad - \frac{1}{2} (\gamma^a)_{\gamma\delta} \psi_a \gamma_I^{\gamma} (\lambda^{\delta J} \mathbf{g}_{-J}^I + \lambda^{\delta J} g_{+J}^I) \\ &\quad \left. + \frac{i}{2} \varepsilon^{abc} (\gamma_a)_{\gamma\delta} \psi_b \gamma_K^{\gamma} \psi_{cL}^{\delta} g_+^{KP} \mathbf{g}_{-L}^P \right) , \end{aligned} \quad (4.9)$$

where the bolded component fields correspond to those of the composite vector multiplet,

$$\mathbf{g}_{-}^{IJ} = \mathbf{G}_{-}^{IJ} | , \quad \lambda_{\alpha}^I = \frac{2}{3} \nabla_{\alpha J} \mathbf{G}_{-}^{IJ} | , \quad \hat{\mathbf{h}}_+^{IJ} = \frac{i}{3} \nabla^{[I} \nabla_{\gamma K} \mathbf{G}_{-}^{J]K} | + 2w \mathbf{g}_{-}^{IJ} , \quad (4.10a)$$

$$v_a = e_a^m V_m | = V_a | + \frac{1}{2} \psi_a \gamma_I^{\alpha} V_{\alpha}^I | , \quad (4.10b)$$

$$\mathbf{f}_{(+ab)} = -\frac{i}{24} \varepsilon_{abc} (\gamma^c)^{\alpha\beta} \nabla_{\alpha}^K \nabla_{\beta}^L \mathbf{G}_{-KL} | - \frac{1}{2} (\psi_{[a}^K \gamma_{b]} \lambda_K) + \frac{i}{2} \psi_a \gamma^K \psi_b \gamma_L \mathbf{g}_{-KL} . \quad (4.10c)$$

The component anti-self dual vector multiplet action is

$$\begin{aligned} S_{\text{VM}}^{(-)} &= - \int d^3x e \left( \varepsilon^{abc} v_{(-a)} \mathbf{f}_{(-b)c} + \frac{1}{4} \hat{\mathbf{h}}_+^{IJ} g_{+IJ} + \frac{1}{4} \hat{\mathbf{h}}_-^{IJ} g_{-IJ} - \frac{i}{2} \lambda^{\alpha I} \lambda_{\alpha I} \right. \\ &\quad - \frac{1}{2} (\gamma^a)_{\gamma\delta} \psi_a \gamma_I^{\gamma} (\lambda^{\delta J} g_{-J}^I + \lambda^{\delta J} \mathbf{g}_{+J}^I) \\ &\quad \left. + \frac{i}{2} \varepsilon^{abc} (\gamma_a)_{\gamma\delta} \psi_b \gamma_K^{\gamma} \psi_{cL}^{\delta} \mathbf{g}_{+}^{KP} g_{-L}^P \right) , \end{aligned} \quad (4.11)$$

where

$$\mathbf{g}_{+}^{IJ} = \mathbf{G}_{+}^{IJ} | , \quad \lambda_{\alpha}^I = \frac{2}{3} \nabla_{\alpha J} \mathbf{G}_{+}^{IJ} | , \quad \hat{\mathbf{h}}_-^{IJ} = \frac{i}{3} \nabla^{[I} \nabla_{\gamma K} \mathbf{G}_{+}^{J]K} | - 2w \mathbf{g}_{+}^{IJ} , \quad (4.12a)$$

$$v_a = e_a{}^m V_m| = V_a| + \frac{1}{2} \psi_a{}^\alpha V_\alpha^I , \quad (4.12b)$$

$$\mathbf{f}_{(-)ab} = -\frac{i}{24} \varepsilon_{abc} (\gamma^c)^{\alpha\beta} \nabla_\alpha^K \nabla_\beta^L \mathbf{G}_{+KL}| - \frac{1}{2} (\psi_{[a}{}^K \gamma_{b]} \boldsymbol{\lambda}_K) + \frac{i}{2} \psi_a{}^{\gamma K} \psi_{b\gamma}{}^L \mathbf{g}_{+KL} . \quad (4.12c)$$

Plugging in the superspace expressions for  $\mathbf{G}_\pm^{IJ}$  one can construct the component fields of the composite vector multiplets. The component fields are found to be

$$\mathbf{g}_\pm^{IJ} = \frac{1}{g_\pm} \hat{h}_\pm^{IJ} - \frac{i}{2g_\pm^3} \lambda_{\pm K}{}^\alpha \Lambda_{\pm\alpha}^{[I} g_\pm^{J]K} \pm \frac{i}{4g_\pm^3} \varepsilon^{IJLP} \lambda_{\pm K}{}^\alpha \lambda_{\pm\alpha L} g_{\pm P}{}^K , \quad (4.13a)$$

$$\begin{aligned} \Lambda_{(\pm)\alpha}^I &= \frac{2}{g_\pm} \nabla_\alpha{}^\gamma \lambda_{(\pm)\gamma}^I + \frac{2}{g_\pm^3} f_{\pm\alpha\beta} \lambda_{(\pm)\gamma}^\beta g_\pm^{IJ} + \frac{1}{3g_\pm^3} \hat{h}_{\mp JK} \lambda_{(\pm)\alpha}^I g_\pm^{JK} \\ &\quad + \frac{2}{3g_\pm^3} \hat{h}_{\mp JK} \lambda_{(\pm)\alpha}^J g_\pm^{KI} + \frac{1}{3g_\pm^2} \hat{h}_{\mp}^{IJ} \lambda_{(\pm)\alpha}^K g_{\pm JK} \\ &\quad + \frac{2}{3g_\pm^3} \nabla_{\alpha\beta} g_{\pm JK} \lambda_{(\pm)}^{\beta I} g_\pm^{JK} + \frac{4}{3g_\pm^3} \nabla_{\alpha\beta} g_{\pm JK} \lambda_{(\pm)}^{\beta J} g_\pm^{KI} \\ &\quad - \frac{2}{g_{(\pm)}^3} \nabla_{\alpha\beta} g_\pm^{IJ} \lambda_{(\pm)}^{\beta K} g_{\pm JK} \\ &\quad \pm \frac{1}{g_\pm} w \lambda_{(\pm)\alpha}^I \pm \frac{8i}{g_\pm} w_{\alpha J} g_\pm^{IJ} + \mathcal{O}(\lambda^2) , \end{aligned} \quad (4.13b)$$

$$\begin{aligned} \hat{\mathbf{h}}_\pm^{IJ} &= \frac{4}{g_\pm} \square g_\pm^{IJ} + \frac{2}{g_\pm^3} f_{\pm ab} f_\pm{}^{ab} g_\pm^{IJ} + \frac{4}{g_\pm^3} \varepsilon^{abc} f_{\pm ab} \nabla_c g_{\pm K}^{[I} g_\pm^{J]K} \\ &\quad - \frac{1}{4g_\pm^3} \hat{h}_{\mp}^{KL} \hat{h}_{\mp KL} g_\pm^{IJ} - \frac{2}{g_\pm^3} g_\pm^{KL} \nabla_a g_{\pm KL} \nabla^a g_\pm^{IJ} \\ &\quad + \frac{1}{g_\pm^3} g_\pm^{IJ} \nabla_a g_{\pm KL} \nabla^a g_\pm^{KL} - \frac{2}{g_\pm} w^2 g_\pm^{IJ} \pm \frac{2}{g_\pm} y g_\pm^{IJ} \\ &\quad + \text{fermion terms} , \end{aligned} \quad (4.13c)$$

$$\begin{aligned} \mathbf{f}_{(\pm)mn} &:= e_m{}^a e_n{}^b \mathbf{f}_{(\pm)ab} \\ &= \partial_{[m} \left( \frac{4}{g_\pm} f_{(\pm)n]} - \frac{2}{g_\pm} V_n{}^{IJ} g_{\pm IJ} \right) - \frac{1}{g_\pm^3} \partial_{[m} g_\pm^{IK} \partial_{n]} g_\pm{}^J{}_K g_{\pm IJ} \\ &\quad + \text{fermion terms} , \end{aligned} \quad (4.13d)$$

where

$$f_{(\pm)}{}^m = \frac{1}{2} \varepsilon^{mnp} f_{(\pm)np} . \quad (4.14)$$

Here we have introduced the following:

$$\nabla_a g_\pm^{IJ} := \mathcal{D}_a g_\pm^{IJ} + \frac{1}{2} \psi_a{}^\alpha [I \lambda_{\pm\alpha}^{J]} \pm \frac{1}{4} \varepsilon^{IJKL} \psi_{aK}{}^\alpha \lambda_{\pm\alpha L} , \quad (4.15a)$$

$$\square g_\pm^{IJ} := \mathcal{D}^a \mathcal{D}_a g_\pm^{IJ} + \frac{1}{4} \mathcal{R} g_\pm^{IJ} + \text{fermion terms} , \quad (4.15b)$$

and<sup>13</sup>

$$\mathcal{D}_a := e_a^m \partial_m - \frac{1}{2} \omega_a^{bc} M_{bc} - \frac{1}{2} V_a^{IJ} N_{IJ} - b_a \mathbb{D} . \quad (4.16)$$

### 4.3 $\mathcal{N} = 4$ topologically massive supergravity in components

To simplify our results it is useful to make use of the gauge freedom to impose some gauge condition. One can always choose a gauge condition where

$$B_A = 0 , \quad G_{\pm} = 1 . \quad (4.17)$$

At the component level these require

$$g_{\pm} = 1 , \quad \lambda_{\alpha}^I = 0 , \quad b_m = 0 . \quad (4.18)$$

The first gauge condition fixes the dilatation transformations, the second fixes the  $S$ -supersymmetry transformations and the third fixes the conformal boosts. For a right  $G^{ij}$  and left  $\bar{G}^{\bar{i}\bar{j}}$  vector multiplet we can use the respective  $SU(2)$  symmetry to fix their lowest components to a constant. This then gives

$$\nabla_a g_{\pm}^{IJ} = 2V_a^{K[I} g_{\pm K}^{J]} . \quad (4.19)$$

With the above gauge conditons we find

$$\begin{aligned} \hat{h}_{\pm}^{IJ} g_{\pm IJ} &= 2\mathcal{R} + 4f_{\pm ab} f_{\pm}^{ab} - \frac{1}{2} \hat{h}_{\mp}^{IJ} \hat{h}_{\mp IJ} \\ &\quad - 2V_{aKL} V^{aKL} + 4V_a^{IK} V^{aJL} g_{\pm IJ} g_{\pm KL} \\ &\quad - 4w^2 \pm 4y + \text{fermion terms} , \end{aligned} \quad (4.20a)$$

$$\hat{h}_{\pm}^{IJ} g_{\pm IJ} = \hat{h}_{\pm}^{IJ} \hat{h}_{\pm IJ} , \quad (4.20b)$$

$$\mathbf{f}_{(\pm)mn} = \partial_{[m} \left( 4f_{(\pm)n]} - 2V_{n]}^{IJ} g_{\pm IJ} \right) + \text{fermion terms} . \quad (4.20c)$$

Using the above conditions one finds (upon integrating by parts) the self-dual vector multiplet action is

$$\begin{aligned} S_{\text{VM}}^{(+)} &= - \int d^3x e \left( \frac{1}{2} \mathcal{R} - f_{(+ab} f_{(+)}^{ab} - 2f_{(+)}^a V_a^{IJ} g_{+IJ} - \frac{1}{2} V_{aKL} V^{aKL} \right. \\ &\quad \left. + V_a^{IK} V^{aJL} g_{+IJ} g_{+KL} + \frac{1}{8} \hat{h}_{-}^{IJ} \hat{h}_{-IJ} - w^2 + y + \text{fermion terms} \right) , \end{aligned} \quad (4.21)$$

---

<sup>13</sup>We have denoted the component vector derivative  $\mathcal{D}_a$  in the same way as the  $SU(2)$  superspace covariant derivative. It should be clear from context to which we are referring to.

while the anti-self-dual vector multiplet action is

$$S_{\text{VM}}^{(-)} = - \int d^3x e \left( \frac{1}{2} \mathcal{R} - f_{(-)ab} f_{(-)}^{ab} - 2f_{(-)}^a V_a^{IJ} g_{-IJ} - \frac{1}{2} V_{aKL} V^{aKL} \right. \\ \left. + V_a^{IK} V^{aJL} g_{-IJ} g_{-KL} + \frac{1}{8} \hat{h}_+^{IJ} \hat{h}_{+IJ} - w^2 - y + \text{fermion terms} \right) . \quad (4.22)$$

The complete component action for minimal  $\mathcal{N} = 4$  topologically massive supergravity (3.1) is then given by

$$\kappa S_{\pm} = \frac{1}{\mu} S_{\text{CSG}} + S_{\text{VM}}^{(\pm)} , \quad (4.23)$$

where  $S_{\text{CSG}}$  is the component action (4.2). As a simple check one can readily verify that the equation of motion on the field  $y$  gives

$$w = \mp \mu , \quad (4.24)$$

which is consistent with the supergravity equation of motion being  $W = \mp \mu G_{\pm}$  in the presence of the vector multiplet compensator.

For completeness we will also give the component action in isospinor notation. The  $\mathcal{N} = 4$  conformal supergravity action (4.2) becomes

$$S_{\text{CSG}} = \frac{1}{8} \int d^3x e \left\{ \varepsilon^{abc} \left( \omega_a^{fg} \mathcal{R}_{bcfg} - \frac{2}{3} \omega_a f^g \omega_{bg}^h \omega_{ch}^f \right. \right. \\ \left. - 4 \mathcal{R}_{ab}^{ij} V_{cij} - \frac{8}{3} V_{ai}^j V_{bj}^k V_{ck}^i \right. \\ \left. - 4 \mathcal{R}_{ab}^{\bar{i}\bar{j}} V_{c\bar{i}\bar{j}} - \frac{8}{3} V_{a\bar{i}}^{\bar{j}} V_{b\bar{j}}^{\bar{k}} V_{c\bar{k}}^{\bar{i}} \right) \\ \left. - 32i w_{ii}^{\alpha} w_{\alpha}^{\bar{i}\bar{i}} - 8wy - 16i \psi_{a\bar{i}\bar{i}}^{\alpha} (\gamma^a)_{\alpha}^{\beta} w_{\beta}^{\bar{i}\bar{i}} w - 2i \varepsilon^{abc} (\gamma_a)_{\alpha\beta} \psi_{b\bar{i}\bar{i}}^{\alpha} \psi_c^{\beta\bar{i}\bar{i}} w^2 \right\} , \quad (4.25)$$

where the component  $\text{SU}(2)$  curvatures  $\mathcal{R}_{ab}^{ij}$  and  $\mathcal{R}_{ab}^{\bar{i}\bar{j}}$  are

$$\mathcal{R}_{ab}^{ij} := 2e_a^m e_b^n \partial_{[m} V_{n]}^{ij} - 2V_{[a}^{ik} V_{b]k}^j , \quad (4.26a)$$

$$\mathcal{R}_{ab}^{\bar{i}\bar{j}} := 2e_a^m e_b^n \partial_{[m} V_{n]}^{\bar{i}\bar{j}} - 2V_{[a}^{\bar{i}\bar{k}} V_{b]\bar{k}}^{\bar{j}} . \quad (4.26b)$$

The self-dual vector multiplet action in isospinor notation is

$$S_{\text{VM}}^{(+)} = - \int d^3x e \left( \frac{1}{2} \mathcal{R} - f_{(+)ab} f_{(+)}^{ab} - 4f_{(+)}^a V_a^{\bar{i}\bar{j}} g_{+\bar{i}\bar{j}} - V_{a\bar{i}\bar{j}} V^{a\bar{i}\bar{j}} \right. \\ \left. + 2V_a^{\bar{i}\bar{k}} V^{a\bar{j}\bar{l}} g_{+\bar{i}\bar{j}} g_{+\bar{k}\bar{l}} + \frac{1}{4} \hat{h}_-^{ij} \hat{h}_{-ij} - w^2 + y + \text{fermion terms} \right) , \quad (4.27)$$

while the anti-self-dual vector multiplet action is

$$S_{\text{VM}}^{(-)} = - \int d^3x e \left( \frac{1}{2} \mathcal{R} - f_{(-)ab} f_{(-)}^{ab} - 4f_{(-)}^a V_a^{ij} g_{-ij} - V_{a ij} V^{a ij} \right)$$



$$+ 2V_a^{ik}V^{ajl}g_{-ij}g_{-kl} + \frac{1}{4}\hat{h}_+^{\bar{i}\bar{j}}\hat{h}_{+\bar{i}\bar{j}} - w^2 - y + \text{fermion terms} \Big) . \quad (4.28)$$

Having derived the component actions for minimal  $\mathcal{N} = 4$  topologically massive supergravity, it is worth elaborating on these results further. For instance, if we consider just one of the vector multiplet actions without the conformal supergravity action, one can see that the equation of motion for  $y$  leads to an inconsistency. This is equivalent to the fact that the superfield equations of motion for the  $\mathcal{N} = 4$  gravitational superfield<sup>14</sup> derived from the actions  $S_{\text{VM}}^{(+)}$  and  $S_{\text{VM}}^{(-)}$  are  $G_+ = 0$  and  $G_- = 0$ , respectively, and these equations are inconsistent with the requirements  $G_{\pm} \neq 0$ . However, one gets consistent equations of motion if one adds the left and right vector multiplets [12] and considers the action

$$S = S_{\text{VM}}^{(+)} + S_{\text{VM}}^{(-)} . \quad (4.29)$$

Now the superfield equation of motion for the  $\mathcal{N} = 4$  gravitational superfield is [12]

$$G_+ - G_- = 0 , \quad (4.30)$$

which is completely consistent. Moreover, this equation is consistent with our gauge conditions because imposing the gauge  $G_+ = 1$  implies  $G_- = 1$ , which in turn implies that the auxiliary field  $y$  cancels. Furthermore, the fields  $w$  and  $\hat{h}^{IJ}$  become auxiliary and their equation of motion is the requirement that they vanish. The equations of motion on the  $\text{SU}(2)$  connections requires  $f_{(-)a} = f_{(+)a} = 0$  and we are left with just the  $\mathcal{N} = 4$  Poincaré supergravity action (up to a normalisation factor)

$$S = - \int d^3x e \mathcal{R} + \text{fermion terms} . \quad (4.31)$$

In the presence of the conformal supergravity action the gauge conditions  $G_+ = G_- = 1$  are no longer consistent [12] and instead one has to use the results in subsection 4.2 in the general gauge. If one also adds to (4.29) the supersymmetric cosmological term [14], the resulting theory corresponds to (2,2) AdS supergravity as was described in detail in [12, 14].

It is worth mentioning some simplifications that can be made to the  $\mathcal{N} = 4$  topologically massive supergravity actions upon using the equations of motion. To illustrate this let us consider the theory with a self-dual vector multiplet. In this case the equation of motion for the  $\text{SU}(2)_\text{L}$  gauge field is

$$\mathcal{R}_{ab}{}^{ij} = 0 , \quad (4.32)$$

---

<sup>14</sup>The  $\mathcal{N} = 4$  gravitational superfield is a scalar prepotential describing the multiplet of  $\mathcal{N} = 4$  conformal supergravity. It is the 3D  $\mathcal{N} = 4$  counterpart of the  $\mathcal{N} = 2$  gravitational superfield in four dimensions [46].

which tells us that the  $SU(2)_L$  gauge field can be completely gauged away. The equation of motion for the auxiliary field  $\hat{h}^{ij}$  sets the auxiliary field to zero and removes it from the action. The equation of motion on  $y$  just sets  $w = -\mu$  and gives rise to a cosmological term. The resulting action is

$$\begin{aligned} \kappa S_+ = \int d^3x e \Big[ & \frac{1}{8\mu} \left\{ \varepsilon^{abc} (\omega_a^{fg} \mathcal{R}_{bcfg} - \frac{2}{3} \omega_{af}^g \omega_{bg}^h \omega_{ch}^f \right. \\ & - 4 \mathcal{R}_{ab}^{\bar{i}\bar{j}} V_{c\bar{i}\bar{j}} - \frac{8}{3} V_{a\bar{i}}^{\bar{j}} V_{b\bar{j}}^{\bar{k}} V_{c\bar{k}}^{\bar{i}}) \Big\} \\ & - \frac{1}{2} \mathcal{R} + \mu^2 + f_{(+)\bar{a}b} f_{(+)}^{ab} + 4 f_{(+)}^a V_a^{\bar{i}\bar{j}} g_{+\bar{i}\bar{j}} + V_{a\bar{i}\bar{j}} V^{a\bar{i}\bar{j}} \\ & \left. - 2 V_a^{\bar{i}\bar{k}} V^{a\bar{j}\bar{l}} g_{+\bar{i}\bar{j}} g_{+\bar{k}\bar{l}} + \text{fermion terms} \right] , \end{aligned} \quad (4.33)$$

Similar simplifications can be made for the anti-self dual vector multiplet action.

We can now show how to derive the supergravity action postulated in [20] from our theory  $S_-$ . The crucial observation is that the  $U(1)$  gauge field appears in the action (4.28) only via its field strength  $f_{(-)ab}$ , and therefore it may be dualised into a scalar field. To implement this, we replace (4.28) with an equivalent first-order action

$$\begin{aligned} S_{\text{FO}}^{(-)} = - \int d^3x e \Big( & \frac{1}{2} \mathcal{R} - f_{(-)ab} f_{(-)}^{ab} - 4 f_{(-)}^a V_a^{ij} g_{ij} - V_{a\bar{i}\bar{j}} V^{a\bar{i}\bar{j}} + 2 V_a^{ik} V^{a\bar{j}\bar{l}} g_{+\bar{i}\bar{j}} g_{+kl} \\ & + \frac{1}{4} \hat{h}_+^{\bar{i}\bar{j}} \hat{h}_{+\bar{i}\bar{j}} - w^2 - y + 2 f_{(-)}^a \mathcal{D}_a \varphi + \text{fermion terms} \Big) , \end{aligned} \quad (4.34)$$

where  $f_{(-)ab}$  is an unconstrained antisymmetric tensor field, and  $\varphi$  a Lagrange multiplier. Varying  $\varphi$  gives  $\mathcal{D}_a f_{(-)}^a = 0$ , and therefore  $f_{(-)ab}$  becomes the field strength of a  $U(1)$  vector multiplet. Then  $S_{\text{FO}}^{(-)}$  turns into the original action (4.28). On the other hand, we may integrate out  $f_{(-)ab}$  from  $S_{\text{FO}}^{(-)}$  using its equation of motion

$$f_{(-)a} = V_a^{ij} g_{ij} - \frac{1}{2} \mathcal{D}_a \varphi . \quad (4.35)$$

Plugging this back into (4.34) gives the dual action

$$\begin{aligned} S_{\text{hyper}}^{(-)} = - \int d^3x e \Big( & \frac{1}{2} \mathcal{R} - \frac{1}{2} \mathcal{D}^a \varphi \mathcal{D}_a \varphi + 2 \mathcal{D}_a \varphi V^{a\bar{i}\bar{j}} g_{+\bar{i}\bar{j}} - 2 V_{a\bar{i}\bar{j}} V^{a\bar{i}\bar{j}} \\ & + \frac{1}{4} \hat{h}_+^{\bar{i}\bar{j}} \hat{h}_{+\bar{i}\bar{j}} - w^2 - y + \text{fermion terms} \Big) , \end{aligned} \quad (4.36)$$

where we used

$$V_a^{ik} V^{a\bar{j}\bar{l}} g_{ij} g_{kl} = V_a^{ij} V^{akl} g_{ij} g_{kl} - \frac{1}{2} V_{a\bar{i}\bar{j}} V^{a\bar{i}\bar{j}} . \quad (4.37)$$

If we impose a Weyl gauge  $\varphi = 1$  and make use of the equation of motion for the auxiliary field  $\hat{h}_+^{\bar{i}\bar{j}}$ , which is  $\hat{h}_+^{\bar{i}\bar{j}} = 0$ , we recover the bosonic matter sector of the topologically massive supergravity action in [20] up to conventions and fermion terms. Since the auxiliary field  $\hat{h}_+^{\bar{i}\bar{j}}$  has been integrated out, the action given in [20] does not appear to be off-shell.

## 5 Discussion

In this paper we constructed minimal topologically massive  $\mathcal{N} = 4$  supergravity. It has several unique features that we summarise here.

- Unlike the other  $\mathcal{N}$ -extended TMSG theories with  $\mathcal{N} \leq 4$  [9, 10, 11, 12], its action cannot be viewed as the supergravity action (with or without a supersymmetric cosmological term) augmented by the conformal supergravity action playing the role of a topological mass term. The point is that the theory becomes inconsistent upon removing the conformal supergravity action, as was explained in section 4.3.
- Our theory makes use of a single superconformal compensator. We recall that all known Poincaré or AdS supergravity theories with eight supercharges in diverse dimensions require, in general, two such compensators in order for the corresponding dynamics to be consistent. One known exception is the off-shell formulation for 4D  $\mathcal{N} = 2$  AdS supergravity given in [47], which makes use of a single massive tensor compensator (described by an unconstrained chiral scalar prepotential) and no compensating vector multiplet.<sup>15</sup> In the case of higher derivative theories, two compensators are no longer required. This was observed in four dimensions for models involving the  $\mathcal{N} = 2$  supersymmetric  $R^2$  term [55], and in three dimensions for topologically massive  $\mathcal{N} = 4$  supergravity [20].
- Our minimal TMSG theory does not allow any supersymmetric cosmological term. However, a cosmological term gets generated at the component level upon integrating out the auxiliary fields. This is manifested in the fact that the critical (4,0) AdS superspace [25] is a maximally supersymmetric solution of the theory.
- The theory has only one coupling constant.
- Our minimal TMSG theory is the first off-shell  $\mathcal{N} = 4$  supergravity theory in three dimensions with the property that the critical (4,0) AdS superspace [25] is a solution of the theory. Upon integrating out the auxiliary fields we recover the model discussed in [20].

---

<sup>15</sup>The vector multiplet has been eaten up by the tensor multiplet. The vector compensator acts as a Stückelberg field to give mass to the tensor multiplet. This is an example of the phenomenon observed originally in [48] and studied in detail in [29, 49, 50, 51, 52, 53, 54].

Both models for minimal topologically massive  $\mathcal{N} = 4$  supergravity constructed in this paper possess dual formulations. They are obtained by replacing the vector multiplet actions  $S_{\text{VM}}^{(+)}$  and  $S_{\text{VM}}^{(-)}$  with off-shell hypermultiplet actions  $S_{\text{HM}}^{(+)}$  and  $S_{\text{HM}}^{(-)}$ , respectively, such that

$$S_{\text{HM}}^{(+)} := -\frac{i}{2\pi} \oint (v_{\text{R}}, dv_{\text{R}}) \int d^{3|8} z E C_{\text{R}}^{(-4)} \Upsilon_{\text{R}}^{(1)} \check{\Upsilon}_{\text{R}}^{(1)} , \quad (5.1)$$

and similarly for the left hypermultiplet action  $S_{\text{HM}}^{(-)}$ . In the dual formulation, its compensating multiplet is the so-called polar hypermultiplet described by the weight-one arctic multiplet  $\Upsilon_{\text{R}}^{(1)}$  and its smile conjugate  $\check{\Upsilon}_{\text{R}}^{(1)}$ . Duality between the theories with actions  $S_{\text{VM}}^{(+)}$  and  $S_{\text{HM}}^{(+)}$  can be shown in complete analogy with the 4D  $\mathcal{N} = 2$  case [29].

### Acknowledgements:

SMK acknowledges the hospitality of the Arnold Sommerfeld Center for Theoretical Physics at the Ludwig Maximilian University of Munich in July 2015, and of the Theoretical Physics Group at Imperial College, London in April 2016. SMK and JN thank the Galileo Galilei Institute for Theoretical Physics for the hospitality and the INFN for partial support during the completion of this work in September 2016. The work of SMK was supported in part by the Australian Research Council, project No. DP140103925. JN acknowledges support from GIF – the German-Israeli Foundation for Scientific Research and Development. I.S. would like to thank DAMTP at Cambridge University for hospitality during the initial stages of this work. I.S. was supported by the DFG Transregional Collaborative Research Centre TRR 33 and the DFG cluster of excellence “Origin and Structure of the Universe”.

## A The geometry of $\mathcal{N} = 4$ conformal superspace

Here we collect the essential details of the  $\mathcal{N} = 4$  superspace geometry of [18]. We refer the reader to [14, 18] for our conventions for 3D spinors.

We begin with a curved three-dimensional  $\mathcal{N} = 4$  superspace  $\mathcal{M}^{3|8}$  parametrized by local bosonic  $(x^m)$  and fermionic coordinates  $(\theta_I^\mu)$ :

$$z^M = (x^m, \theta_I^\mu) , \quad (A.1)$$

where  $m = 0, 1, 2$ ,  $\mu = 1, 2$  and  $I = 1, \dots, 4$ . The structure group is chosen to be  $\text{Osp}(4|4, \mathbb{R})$  and the covariant derivatives are postulated to have the form

$$\nabla_A = E_A - \omega_A{}^b{}_{\underline{b}} X_{\underline{b}} = E_A - \frac{1}{2} \Omega_A{}^{bc} M_{bc} - \frac{1}{2} \Phi_A{}^{PQ} N_{PQ} - B_A \mathbb{D} - \mathfrak{F}_A{}^B K_B . \quad (A.2)$$

Here  $E_A = E_A^M \partial_M$  is the inverse vielbein,  $M_{ab}$  are the Lorentz generators,  $N_{IJ}$  are generators of the  $\text{SO}(4)$  group,  $\mathbb{D}$  is the dilatation generator and  $K_A = (K_a, S_\alpha^I)$  are the special superconformal generators.

The Lorentz generators obey

$$[M_{ab}, M_{cd}] = 2\eta_{c[a} M_{b]d} - 2\eta_{d[a} M_{b]c} , \quad (\text{A.3a})$$

$$[M_{ab}, \nabla_c] = 2\eta_{c[a} \nabla_{b]} , \quad [M_{\alpha\beta}, \nabla_\gamma^I] = \varepsilon_{\gamma(\alpha} \nabla_{\beta)}^I . \quad (\text{A.3b})$$

The  $\text{SO}(4)$  and dilatation generators obey

$$[N_{KL}, N^{IJ}] = 2\delta_{[K}^I N_{L]}^J - 2\delta_{[K}^J N_{L]}^I , \quad [N_{KL}, \nabla_\alpha^I] = 2\delta_{[K}^I \nabla_{\alpha L]} , \quad (\text{A.3c})$$

$$[\mathbb{D}, \nabla_a] = \nabla_a , \quad [\mathbb{D}, \nabla_\alpha^I] = \frac{1}{2} \nabla_\alpha^I . \quad (\text{A.3d})$$

The Lorentz and  $\text{SO}(4)$  generators act on the special conformal generators  $K_A$  as

$$[M_{ab}, K_c] = 2\eta_{c[a} K_{b]} , \quad [M_{\alpha\beta}, S_\gamma^I] = \varepsilon_{\gamma(\alpha} S_{\beta)}^I , \quad (\text{A.3e})$$

$$[N_{KL}, S_\alpha^I] = 2\delta_{[K}^I S_{\alpha L]} , \quad (\text{A.3f})$$

while the dilatation generator acts on  $K_A$  as

$$[\mathbb{D}, K_a] = -K_a , \quad [\mathbb{D}, S_\alpha^I] = -\frac{1}{2} S_\alpha^I . \quad (\text{A.3g})$$

Among themselves, the generators  $K_A$  obey the algebra

$$\{S_\alpha^I, S_\beta^J\} = 2i\delta^{IJ}(\gamma^c)_{\alpha\beta} K_c . \quad (\text{A.3h})$$

Finally, the algebra of  $K_A$  with  $\nabla_A$  is given by

$$[K_a, \nabla_b] = 2\eta_{ab} \mathbb{D} + 2M_{ab} , \quad (\text{A.3i})$$

$$[K_a, \nabla_\alpha^I] = -i(\gamma_a)_\alpha{}^\beta S_\beta^I , \quad (\text{A.3j})$$

$$[S_\alpha^I, \nabla_a] = i(\gamma_a)_\alpha{}^\beta \nabla_\beta^I , \quad (\text{A.3k})$$

$$\{S_\alpha^I, \nabla_\beta^J\} = 2\varepsilon_{\alpha\beta} \delta^{IJ} \mathbb{D} - 2\delta^{IJ} M_{\alpha\beta} - 2\varepsilon_{\alpha\beta} N^{IJ} . \quad (\text{A.3l})$$

The covariant derivatives obey the (anti-)commutation relations of the form

$$\begin{aligned} [\nabla_A, \nabla_B] &= -T_{AB}{}^C \nabla_C - \frac{1}{2} R(M)_{AB}{}^{cd} M_{cd} - \frac{1}{2} R(N)_{AB}{}^{PQ} N_{PQ} \\ &\quad - R(\mathbb{D})_{AB} \mathbb{D} - R(S)_{ABK}{}^\gamma S_\gamma^K - R(K)_{AB}{}^c K_c , \end{aligned} \quad (\text{A.4})$$

where  $T_{AB}{}^C$  is the torsion, and  $R(M)_{AB}{}^{cd}$ ,  $R(N)_{AB}{}^{PQ}$ ,  $R(\mathbb{D})_{AB}$ ,  $R(S)_{ABK}{}^\gamma$  and  $R(K)_{AB}{}^c$  are the curvatures corresponding to the Lorentz,  $\text{SO}(4)$ , dilatation,  $S$ -supersymmetry and special conformal boosts, respectively.

The full gauge group of conformal supergravity,  $\mathcal{G}$ , is generated by *covariant general coordinate transformations*,  $\delta_{\text{cgct}}$ , associated with a parameter  $\xi^A$  and *standard superconformal transformations*,  $\delta_{\mathcal{H}}$ , associated with a parameter  $\Lambda^a$ . The latter include the dilatation, Lorentz,  $\text{SO}(4)$ , and special conformal (bosonic and fermionic) transformations. The covariant derivatives transform as

$$\delta_{\mathcal{G}} \nabla_A = [\mathcal{K}, \nabla_A] , \quad (\text{A.5})$$

where  $\mathcal{K}$  denotes the first-order differential operator

$$\mathcal{K} = \xi^C \nabla_C + \frac{1}{2} \Lambda^{ab} M_{ab} + \frac{1}{2} \Lambda^{IJ} N_{IJ} + \Lambda \mathbb{D} + \Lambda^A K_A . \quad (\text{A.6})$$

Covariant (or tensor) superfields transform as

$$\delta_{\mathcal{G}} T = \mathcal{K} T . \quad (\text{A.7})$$

In order to describe the Weyl multiplet of conformal supergravity, some of the components of the torsion and curvatures must be constrained. Following [18], the spinor derivative torsion and curvatures are chosen to resemble super-Yang Mills

$$\{\nabla_{\alpha}^I, \nabla_{\beta}^J\} = -2i\varepsilon_{\alpha\beta} \mathcal{W}^{IJ} , \quad (\text{A.8})$$

where  $\mathcal{W}^{IJ}$  is some operator that takes values in the superconformal algebra, with  $P_A$  replaced by  $\nabla_A$ . In [18] it was shown how to constrain  $\mathcal{W}^{IJ}$  entirely in terms of the super Cotton tensor (or scalar for  $\mathcal{N} = 4$ ). The super Cotton scalar  $W$ , is a primary superfield of dimension-1,

$$S_{\alpha}^I W = 0 , \quad K_a W = 0 , \quad \mathbb{D} W = W . \quad (\text{A.9})$$

The algebra of covariant derivatives is

$$\begin{aligned} \{\nabla_{\alpha}^I, \nabla_{\beta}^J\} &= 2i\delta^{IJ} \nabla_{\alpha\beta} + i\varepsilon_{\alpha\beta} \varepsilon^{IJKL} W N_{KL} - i\varepsilon_{\alpha\beta} \varepsilon^{IJKL} (\nabla_K^{\gamma} W) S_{\gamma L} \\ &\quad + \frac{1}{4} \varepsilon_{\alpha\beta} (\gamma^c)^{\gamma\delta} \varepsilon^{IJKL} (\nabla_{\gamma K} \nabla_{\delta L} W) K_c , \end{aligned} \quad (\text{A.10a})$$

$$\begin{aligned} [\nabla_a, \nabla_{\beta}^J] &= \frac{1}{2} \varepsilon^{JPQK} (\gamma_a)_{\beta\gamma} (\nabla_K^{\gamma} W) N_{PQ} \\ &\quad - \frac{1}{4} (\gamma_a)_{\beta\gamma} \varepsilon^{JKLP} (\nabla_L^{\gamma} \nabla_P^{\delta} W) S_{\delta K} \\ &\quad - \frac{i}{24} (\gamma_a)_{\beta\gamma} (\gamma^c)_{\delta\rho} \varepsilon^{JKLP} (\nabla_K^{\gamma} \nabla_L^{\delta} \nabla_P^{\rho} W) K_c , \end{aligned} \quad (\text{A.10b})$$

$$\begin{aligned} [\nabla_a, \nabla_b] &= \frac{1}{8} \varepsilon_{abc} (\gamma^c)_{\alpha\beta} \varepsilon^{PQIJ} \left( i (\nabla_I^{\alpha} \nabla_J^{\beta} W) N_{PQ} \right. \\ &\quad \left. + \frac{i}{3} \varepsilon^{LIJK} (\nabla_I^{\alpha} \nabla_J^{\beta} \nabla_K^{\gamma} W) S_{\gamma L} \right) \end{aligned}$$

$$+ \frac{1}{24}(\gamma^d)_{\gamma\delta}\varepsilon^{IJKL}(\nabla_I^\alpha\nabla_J^\beta\nabla_K^\gamma\nabla_L^\delta W)K_d) , \quad (\text{A.10c})$$

where the super Cotton scalar  $W$  satisfies the following dimension 2 Bianchi identity

$$\nabla^{\alpha I}\nabla_\alpha^J W = \frac{1}{4}\delta^{IJ}\nabla_P^\alpha\nabla_\alpha^P W . \quad (\text{A.11})$$

For each  $\text{SO}(4)$  vector  $V_I$  we can associate a second-rank isospinor  $V_{i\bar{i}}$

$$V_I \leftrightarrow V_{i\bar{i}} := (\tau^I)_{i\bar{i}}V_{i\bar{i}} , \quad (V_{i\bar{i}})^* = V^{\bar{i}i} . \quad (\text{A.12})$$

The original  $\text{SO}(4)$  connection turns into a sum of two  $\text{SU}(2)$  connections

$$\Phi_A = (\Phi_L)_A + (\Phi_R)_A , \quad (\Phi_L)_A = \Phi_A^{kl}L_{kl} , \quad (\Phi_R)_A = \Phi_A^{\bar{k}\bar{l}}R_{\bar{k}\bar{l}} . \quad (\text{A.13})$$

Here  $L_{kl}$  is the  $\text{SU}(2)_L$  generator and  $R_{\bar{k}\bar{l}}$  is the  $\text{SU}(2)_R$  generator. They are related to the  $\text{SO}(4)$  generators  $N_{KL}$  as

$$N_{KL} \rightarrow N_{k\bar{k}l\bar{l}} = \varepsilon_{\bar{k}l}L_{kl} + \varepsilon_{kl}R_{\bar{k}\bar{l}} . \quad (\text{A.14})$$

The left and right operators act on the covariant derivatives as

$$[L^{kl}, \nabla_\alpha^{\bar{i}i}] = \varepsilon^{i(k}\nabla_\alpha^{l)\bar{i}} , \quad [R^{k\bar{l}}, \nabla_\alpha^{\bar{i}i}] = \varepsilon^{\bar{i}(k}\nabla_\alpha^{\bar{l})i} . \quad (\text{A.15})$$

In the isospinor notation, the Bianchi identity on  $W$  becomes

$$\nabla^{\alpha i\bar{i}}\nabla_\alpha^{j\bar{j}}W = \frac{1}{4}\varepsilon^{ij}\varepsilon^{\bar{i}\bar{j}}\nabla_{k\bar{k}}^\alpha\nabla_\alpha^{k\bar{k}}W . \quad (\text{A.16})$$

The algebra of spinor covariant derivatives becomes

$$\begin{aligned} \{\nabla_\alpha^{\bar{i}i}, \nabla_\beta^{j\bar{j}}\} &= 2i\varepsilon^{ij}\varepsilon^{\bar{i}\bar{j}}\nabla_{\alpha\beta} + 2i\varepsilon_{\alpha\beta}\varepsilon^{\bar{i}\bar{j}}WL^{ij} - 2i\varepsilon_{\alpha\beta}\varepsilon^{ij}WR^{\bar{i}\bar{j}} \\ &\quad - i\varepsilon_{\alpha\beta}\varepsilon^{ij}\nabla_\gamma^{\bar{i}i}W S_\gamma^{k\bar{j}} + i\varepsilon_{\alpha\beta}\varepsilon^{\bar{i}\bar{j}}\nabla_\gamma^{j\bar{k}}W S_\gamma^{i\bar{l}} \\ &\quad + \frac{1}{4}\varepsilon_{\alpha\beta}\left(\varepsilon^{ij}\nabla_{\gamma k}^{\bar{i}i}\nabla_\delta^{k\bar{j}}W - \varepsilon^{\bar{i}\bar{j}}\nabla_{\gamma \bar{k}}^j\nabla_\delta^{\bar{k}\bar{l}}W\right)K^{\gamma\delta} \end{aligned} \quad (\text{A.17})$$

and the action of the  $S$ -supersymmetry generator on  $\nabla_\alpha^{\bar{i}i}$  is

$$\{S_\alpha^{\bar{i}i}, \nabla_\beta^{j\bar{j}}\} = 2\varepsilon_{\alpha\beta}\varepsilon^{ij}\varepsilon^{\bar{i}\bar{j}}\mathbb{D} - 2\varepsilon^{ij}\varepsilon^{\bar{i}\bar{j}}M_{\alpha\beta} + 2\varepsilon_{\alpha\beta}\varepsilon^{\bar{i}\bar{j}}L^{ij} + 2\varepsilon_{\alpha\beta}\varepsilon^{ij}R^{\bar{i}\bar{j}} . \quad (\text{A.18})$$

## B The geometry of SO(4) superspace

For many applications it is useful to work with a superspace formulation with a smaller structure group than that of conformal superspace. The superspace formulation of [14, 27], known as SO(4) superspace, provides such a formulation and may be obtained from conformal superspace via a degauging procedure [18]. For the  $\mathcal{N} = 4$  case one chooses the structure group to be SO(4). The SO(4) superspace formulation for  $\mathcal{N} = 4$  conformal supergravity has been used to construct general off-shell supergravity-matter couplings [14].

The covariant derivatives have the form:

$$\mathcal{D}_A = E_A - \Omega_A - \Phi_A . \quad (\text{B.1})$$

Here  $E_A = E_A^M(z)\partial_M$  is the supervielbein, with  $\partial_M = \partial/\partial z^M$ ,  $\Omega_A$  is the Lorentz connection, and  $\Phi_A = \frac{1}{2}\Phi_A^{KL}N_{KL}$  is the SO(4)-connection. The supergravity gauge group is generated by local transformations of the form

$$\delta_K \mathcal{D}_A = [K, \mathcal{D}_A] , \quad K = K^C(z)\mathcal{D}_C + \frac{1}{2}K^{cd}(z)M_{cd} + \frac{1}{2}K^{PQ}(z)N_{PQ} , \quad (\text{B.2})$$

with all the gauge parameters obeying natural reality conditions.

The covariant derivatives satisfy the (anti)commutation relations

$$[\mathcal{D}_A, \mathcal{D}_B] = -T_{AB}^C \mathcal{D}_C - \frac{1}{2}R_{AB}^{KL}N_{KL} - \frac{1}{2}R_{AB}^{cd}M_{cd} , \quad (\text{B.3})$$

with  $T_{AB}^C$  the torsion,  $R_{AB}^{cd}$  the Lorentz curvature and  $R_{AB}^{KL}$  the SO(4) curvature. The algebra of covariant derivatives must be constrained to describe conformal supergravity. The appropriate constraints [27] lead to the following anti-commutation relation [14]:

$$\begin{aligned} \{\mathcal{D}_\alpha^I, \mathcal{D}_\beta^J\} &= 2i\delta^{IJ}(\gamma^c)_{\alpha\beta}\mathcal{D}_c - 2i\varepsilon_{\alpha\beta}C^{\gamma\delta IJ}M_{\gamma\delta} - 4iS^{IJ}M_{\alpha\beta} \\ &\quad + \left(i\varepsilon_{\alpha\beta}W^{IJKL} - 4i\varepsilon_{\alpha\beta}S^{K[I}\delta^{J]L} + iC_{\alpha\beta}^{KL}\delta^{IJ} - 4iC_{\alpha\beta}^{K(I}\delta^{J)L}\right)N_{KL} . \end{aligned} \quad (\text{B.4a})$$

Here the dimension-1 components are real and satisfy the symmetry properties

$$W^{IJKL} = W^{[IJKL]} = \varepsilon^{IJKL}W , \quad S^{IJ} = S^{(IJ)} , \quad C_a^{IJ} = C_a^{[IJ]} . \quad (\text{B.5})$$

It is useful to decompose the torsion superfield  $S^{IJ}$  into its trace ( $\mathcal{S}$ ) and traceless ( $\mathcal{S}^{IJ}$ ) parts as

$$S^{IJ} = \mathcal{S}\delta^{IJ} + \mathcal{S}^{IJ} , \quad \mathcal{S} = \frac{1}{\mathcal{N}}\delta_{IJ}S^{IJ} , \quad \delta_{IJ}\mathcal{S}^{IJ} = 0 . \quad (\text{B.6})$$



The torsion superfields satisfy the Bianchi identities

$$\mathcal{D}_\alpha^I \mathcal{S}^{JK} = 2\mathcal{T}_\alpha^{I(JK)} + \mathcal{S}_\alpha^{(J} \delta^{K)I} - \frac{1}{\mathcal{N}} \mathcal{S}_\alpha^I \delta^{JK} , \quad (\text{B.7a})$$

$$\begin{aligned} \mathcal{D}_\alpha^I C_{\beta\gamma}^{JK} &= \frac{2}{3} \varepsilon_{\alpha(\beta} \left( C_{\gamma)}^{IJK} + 3\mathcal{T}_{\gamma)}^{JKI} + 4(\mathcal{D}_{\gamma)}^{[J} \mathcal{S}) \delta^{K]I} + \frac{(\mathcal{N}-4)}{\mathcal{N}} \mathcal{S}_{\gamma)}^{[J} \delta^{K]I} \right) \\ &\quad + C_{\alpha\beta\gamma}^{IJK} - 2C_{\alpha\beta\gamma}^{[J} \delta^{K]I} , \end{aligned} \quad (\text{B.7b})$$

$$0 = \left( \mathcal{D}^{\gamma(I} \mathcal{D}_{\gamma}^{J)} - \frac{1}{4} \delta^{IJ} \mathcal{D}^{\gamma K} \mathcal{D}_{\gamma K} - 4i\mathcal{S}^{IJ} \right) W . \quad (\text{B.7c})$$

It is often useful to make use of the isomorphism  $\text{SO}(4) \cong (\text{SU}(2)_L \times \text{SU}(2)_R)/\mathbb{Z}_2$  and make use of isospinor notation,  $\mathcal{D}_\alpha^I \rightarrow \mathcal{D}_\alpha^{i\bar{i}}$ , by replacing each  $\text{SO}(4)$  vector index by a pair of isospinor ones. For our notation and conventions we refer the reader to [14].

After introducing isospinor notation, the covariant derivatives are

$$\mathcal{D}_A = (\mathcal{D}_a, \mathcal{D}_\alpha^{i\bar{i}}) = E_A - \Omega_A - \Phi_A , \quad (\text{B.8})$$

where the original  $\text{SO}(4)$  connection  $\Phi_A$  now turns into a sum of two  $\text{SU}(2)$  connections

$$\Phi_A = (\Phi_L)_A + (\Phi_R)_A , \quad (\Phi_L)_A = \Phi_A^{kl} L_{kl} , \quad (\Phi_R)_A = \Phi_A^{\bar{k}\bar{l}} R_{\bar{k}\bar{l}} . \quad (\text{B.9})$$

The two  $\text{SU}(2)$  generators act on the spinor covariant derivatives  $\mathcal{D}_\alpha^{i\bar{i}} := \mathcal{D}_\alpha^I (\tau_I)^{i\bar{i}}$  as follows:

$$[L^{kl}, \mathcal{D}_\alpha^{i\bar{i}}] = \varepsilon^{i(k} \mathcal{D}_\alpha^{l)\bar{i}} , \quad [R^{\bar{k}\bar{l}}, \mathcal{D}_\alpha^{i\bar{i}}] = \varepsilon^{\bar{i}(\bar{k}} \mathcal{D}_\alpha^{i\bar{l}} . \quad (\text{B.10})$$

The algebra of spinor covariant derivatives is

$$\begin{aligned} \{\mathcal{D}_\alpha^{i\bar{i}}, \mathcal{D}_\beta^{j\bar{j}}\} &= 2i\varepsilon^{ij} \varepsilon^{\bar{i}\bar{j}} (\gamma^c)_{\alpha\beta} \mathcal{D}_c + 2i\varepsilon_{\alpha\beta} \varepsilon^{\bar{i}\bar{j}} (2\mathcal{S} + X) L^{ij} - 2i\varepsilon_{\alpha\beta} \varepsilon^{ij} \mathcal{S}^{kl\bar{i}\bar{j}} L_{kl} + 4iC_{\alpha\beta}^{\bar{i}\bar{j}} L^{ij} \\ &\quad + 2i\varepsilon_{\alpha\beta} \varepsilon^{ij} (2\mathcal{S} - X) R^{\bar{i}\bar{j}} - 2i\varepsilon_{\alpha\beta} \varepsilon^{\bar{i}\bar{j}} \mathcal{S}^{ij\bar{k}\bar{l}} R_{\bar{k}\bar{l}} + 4iC_{\alpha\beta}^{ij} R^{\bar{i}\bar{j}} \\ &\quad + 2i\varepsilon_{\alpha\beta} (\varepsilon^{\bar{i}\bar{j}} C^{\gamma\delta ij} + \varepsilon^{ij} C^{\gamma\delta \bar{i}\bar{j}}) M_{\gamma\delta} - 4i(\mathcal{S}^{ij\bar{i}\bar{j}} + \varepsilon^{ij} \varepsilon^{\bar{i}\bar{j}} \mathcal{S}) M_{\alpha\beta} , \end{aligned} \quad (\text{B.11})$$

where the torsion components satisfy certain Bianchi identities given in [14].<sup>16</sup>

## C Super-Weyl gauge conditions

In this appendix we show how one can use the super-Weyl freedom to impose certain gauge conditions in  $\text{SO}(4)$  superspace. In particular, within the  $\text{SO}(4)$  (or  $\text{SU}(2)_L \times \text{SU}(2)_R$ ) superspace formulation we will show that one can impose either

$$C_a^{\bar{i}\bar{j}} = 0 , \quad 2S + W = 0 \quad (\text{C.1})$$

---

<sup>16</sup> As compared to [14], we have relabelled the superfield  $B_{\alpha\beta}^{ij}$  by  $C_{\alpha\beta}^{ij}$ .

or

$$C_a^{ij} = 0, \quad 2S - W = 0. \quad (\text{C.2})$$

We begin by introducing, within the  $\text{SO}(4)$  superspace geometry, an off-shell self-dual vector multiplet  $G^{\bar{i}\bar{j}}$  and an anti-self-dual vector multiplet  $G^{ij}$ . They are constrained by the differential constraints for  $\mathcal{O}(2)$  multiplets

$$\mathcal{D}_\alpha^{i(\bar{i}} G^{\bar{j}\bar{k})} = 0, \quad \mathcal{D}_\alpha^{(\bar{i}\bar{i}} G^{jk)} = 0. \quad (\text{C.3})$$

Using these constraints it is possible to build some of the components of the torsion in terms of these multiplets. In particular, one finds

$$2S - W = \frac{iG_+}{8} \mathcal{D}^{\gamma\bar{i}\bar{i}} \mathcal{D}_{\gamma\bar{i}\bar{i}} G_+^{-1}, \quad (\text{C.4a})$$

$$2S + W = \frac{iG_-}{8} \mathcal{D}^{\gamma\bar{i}\bar{i}} \mathcal{D}_{\gamma\bar{i}\bar{i}} G_-^{-1}, \quad (\text{C.4b})$$

$$C_{\alpha\beta}^{ij} = -\frac{i}{4} G_+ \mathcal{D}_\alpha^{(i\bar{k}} \mathcal{D}_{\beta}^{j)\bar{k}} G_+^{-1}, \quad (\text{C.4c})$$

$$C_{\alpha\beta}^{\bar{i}\bar{j}} = -\frac{i}{4} G_- \mathcal{D}_\alpha^{k(\bar{i}} \mathcal{D}_{\beta k}^{\bar{j})} G_-^{-1}, \quad (\text{C.4d})$$

$$S^{(k}_{\bar{p}} \bar{i}\bar{j} G^{l)p} = -\frac{i}{16} \{ \mathcal{D}^{\gamma p(\bar{i}} \mathcal{D}_{\gamma p}^{\bar{j})} \} G^{kl}, \quad (\text{C.4e})$$

$$S^{ij}_{\bar{p}} (\bar{k} G^{\bar{l})\bar{p}} = -\frac{i}{16} \{ \mathcal{D}^{\gamma(i\bar{p}} \mathcal{D}_{\gamma}^{j)\bar{p}} \} G^{\bar{k}\bar{l}}, \quad (\text{C.4f})$$

where  $G_+^2 = G^{\bar{i}\bar{j}} G_{\bar{i}\bar{j}}$  and  $G_-^2 = G^{ij} G_{ij}$ .

The vector multiplets transform homogeneously under super-Weyl transformations

$$G^{\bar{i}\bar{j}} \rightarrow e^\sigma G^{\bar{i}\bar{j}}, \quad G^{ij} \rightarrow e^\sigma G^{ij}, \quad (\text{C.5})$$

which tells us that the super-Weyl freedom can be completely fixed by imposing the gauge condition  $G_+ = 1$  or  $G_- = 1$ . If we impose  $G_+ = 1$  we find the conditions (C.1), while if we impose  $G_- = 1$  we find the conditions (C.2). Therefore, these conditions can always be imposed by an appropriate super-Weyl transformation.

## References

- [1] W. Siegel, “Unextended superfields in extended supersymmetry,” Nucl. Phys. B **156**, 135 (1979).
- [2] J. F. Schonfeld, “A mass term for three-dimensional gauge fields,” Nucl. Phys. B **185**, 157 (1981).
- [3] S. Deser, R. Jackiw and S. Templeton, “Three-dimensional massive gauge theories,” Phys. Rev. Lett. **48**, 975 (1982).

- [4] S. Deser, R. Jackiw and S. Templeton, “Topologically massive gauge theories,” *Annals Phys.* **140**, 372 (1982) [Erratum-ibid. **185**, 406 (1988)].
- [5] P. van Nieuwenhuizen, “D = 3 conformal supergravity and Chern-Simons terms,” *Phys. Rev. D* **32**, 872 (1985).
- [6] J. H. Horne and E. Witten, “Conformal gravity in three dimensions as a gauge theory,” *Phys. Rev. Lett.* **62**, 501 (1989).
- [7] A. Achúcarro and P. K. Townsend, “A Chern-Simons action for three-dimensional anti-de Sitter supergravity theories,” *Phys. Lett. B* **180**, 89 (1986).
- [8] E. Witten, “(2+1)-dimensional gravity as an exactly soluble system,” *Nucl. Phys. B* **311**, 46 (1988).
- [9] S. Deser and J. H. Kay, “Topologically massive supergravity,” *Phys. Lett. B* **120**, 97 (1983).
- [10] S. Deser, “Cosmological topological supergravity,” in *Quantum Theory Of Gravity*, S. M. Christensen (Ed.), Adam Hilger, Bristol, 1984, pp. 374-381.
- [11] S. M. Kuzenko, U. Lindström, M. Roček, I. Sachs and G. Tartaglino-Mazzucchelli, “Three-dimensional N=2 supergravity theories: From superspace to components,” *Phys. Rev. D* **89**, 085028 (2014) [arXiv:1312.4267 [hep-th]].
- [12] S. M. Kuzenko and J. Novak, “Supergravity-matter actions in three dimensions and Chern-Simons terms,” *JHEP* **1405**, 093 (2014) [arXiv:1401.2307 [hep-th]].
- [13] S. J. Gates Jr., M. T. Grisaru, M. Roček and W. Siegel, *Superspace, or One Thousand and One Lessons in Supersymmetry*, Benjamin/Cummings (Reading, MA), 1983, hep-th/0108200.
- [14] S. M. Kuzenko, U. Lindström and G. Tartaglino-Mazzucchelli, “Off-shell supergravity-matter couplings in three dimensions,” *JHEP* **1103**, 120 (2011) [arXiv:1101.4013 [hep-th]].
- [15] S. M. Kuzenko and G. Tartaglino-Mazzucchelli, “Three-dimensional N=2 (AdS) supergravity and associated supercurrents,” *JHEP* **1112**, 052 (2011) [arXiv:1109.0496 [hep-th]].
- [16] M. Roček and P. van Nieuwenhuizen, “N ≥ 2 supersymmetric Chern-Simons terms as d = 3 extended conformal supergravity,” *Class. Quant. Grav.* **3**, 43 (1986).
- [17] D. Butter, S. M. Kuzenko, J. Novak and G. Tartaglino-Mazzucchelli, “Conformal supergravity in three dimensions: Off-shell actions,” *JHEP* **1310**, 073 (2013) [arXiv:1306.1205 [hep-th]].
- [18] D. Butter, S. M. Kuzenko, J. Novak and G. Tartaglino-Mazzucchelli, “Conformal supergravity in three dimensions: New off-shell formulation,” *JHEP* **1309**, 072 (2013) [arXiv:1305.3132 [hep-th]].
- [19] D. Z. Freedman and A. Van Proeyen, *Supergravity*, Cambridge University Press, Cambridge, 2012.
- [20] F. Lauf and I. Sachs, “On topologically massive gravity with extended supersymmetry,” *Phys. Rev. D* **94**, 065028 (2016) [arXiv:1605.00103 [hep-th]].
- [21] X. Chu and B. E. W. Nilsson, “Three-dimensional topologically gauged N=6 ABJM type theories,” *JHEP* **1006** (2010) 057 [arXiv:0906.1655 [hep-th]].

- [22] U. Gran, J. Greitz, P. S. Howe and B. E. W. Nilsson, “Topologically gauged superconformal Chern-Simons matter theories,” JHEP **1212** (2012) 046 [arXiv:1204.2521 [hep-th]].
- [23] B. E. W. Nilsson, “Critical solutions of topologically gauged  $N = 8$  CFTs in three dimensions,” JHEP **1404** (2014) 107 [arXiv:1304.2270 [hep-th]].
- [24] W. Li, W. Song and A. Strominger, “Chiral gravity in three dimensions,” JHEP **0804**, 082 (2008) [arXiv:0801.4566 [hep-th]].
- [25] S. M. Kuzenko, U. Lindström and G. Tartaglino-Mazzucchelli, “Three-dimensional (p,q) AdS superspaces and matter couplings,” JHEP **1208**, 024 (2012) [arXiv:1205.4622 [hep-th]].
- [26] R. Brooks and S. J. Gates Jr., “Extended supersymmetry and super- $BF$  gauge theories,” Nucl. Phys. B **432**, 205 (1994) [arXiv:hep-th/9407147].
- [27] P. S. Howe, J. M. Izquierdo, G. Papadopoulos and P. K. Townsend, “New supergravities with central charges and Killing spinors in 2+1 dimensions,” Nucl. Phys. B **467**, 183 (1996) [arXiv:hep-th/9505032].
- [28] S. M. Kuzenko and I. B. Samsonov, “Superconformal Chern-Simons-matter theories in  $\mathcal{N} = 4$  superspace,” Phys. Rev. D **92**, no. 10, 105007 (2015) [arXiv:1507.05377 [hep-th]].
- [29] S. M. Kuzenko, “On  $N = 2$  supergravity and projective superspace: Dual formulations,” Nucl. Phys. B **810**, 135 (2009) [arXiv:0807.3381 [hep-th]].
- [30] B. de Wit, R. Philippe and A. Van Proeyen, “The improved tensor multiplet in  $N = 2$  supergravity,” Nucl. Phys. B **219**, 143 (1983).
- [31] S. M. Kuzenko, J. Novak and G. Tartaglino-Mazzucchelli, “ $N=6$  superconformal gravity in three dimensions from superspace,” JHEP **1401**, 121 (2014) [arXiv:1308.5552 [hep-th]].
- [32] S. M. Kuzenko, “Symmetries of curved superspace,” JHEP **1303**, 024 (2013) [arXiv:1212.6179 [hep-th]].
- [33] S. M. Kuzenko, J. Novak and G. Tartaglino-Mazzucchelli, “Symmetries of curved superspace in five dimensions,” JHEP **1410**, 175 (2014) [arXiv:1406.0727 [hep-th]].
- [34] S. M. Kuzenko, “Supersymmetric spacetimes from curved superspace,” PoS CORFU **2014**, 140 (2015) [arXiv:1504.08114 [hep-th]].
- [35] S. M. Kuzenko and G. Tartaglino-Mazzucchelli, “ $N = 4$  supersymmetric Yang-Mills theories in  $AdS_3$ ,” JHEP **1405**, 018 (2014) [arXiv:1402.3961 [hep-th]].
- [36] B. DeWitt, *Supermanifolds*, Cambridge University Press, Cambridge, 1992.
- [37] L. Frappat, A. Sciarrino and P. Sorba, “Dictionary on Lie superalgebras,” hep-th/9607161; *Dictionary on Lie Algebras and Superalgebras*, Academic Press, New York, 2000.
- [38] I. A. Bandos, E. Ivanov, J. Lukierski and D. Sorokin, “On the superconformal flatness of AdS superspaces,” JHEP **0206**, 040 (2002). [hep-th/0205104].
- [39] D. Anninos, W. Li, M. Padi, W. Song and A. Strominger, “Warped AdS(3) black holes,” JHEP **0903**, 130 (2009) [arXiv:0807.3040 [hep-th]].
- [40] N. S. Deger, A. Kaya, H. Samtleben and E. Sezgin, “Supersymmetric warped AdS in extended topologically massive supergravity,” Nucl. Phys. B **884**, 106 (2014) [arXiv:1311.4583 [hep-th]].

- [41] N. S. Deger and G. Moutsopoulos, “Supersymmetric solutions of  $N = (2, 0)$  topologically massive supergravity,” *Class. Quant. Grav.* **33**, no. 15, 155006 (2016) [arXiv:1602.07263 [hep-th]].
- [42] G. Moutsopoulos, “Warped anti-de Sitter in 3d  $(2, 0)$  supergravity,” arXiv:1602.08733 [hep-th].
- [43] N. S. Deger, A. Kaya, H. Samtleben and E. Sezgin, “Supersymmetric warped AdS in extended topologically massive supergravity,” *Nucl. Phys. B* **884**, 106 (2014) [arXiv:1311.4583 [hep-th]].
- [44] S. Deger, A. Kaya, E. Sezgin and P. Sundell, “Spectrum of  $D = 6$ ,  $N=4$  supergravity on AdS in three-dimensions  $\times S^3$ ,” *Nucl. Phys. B* **536**, 110 (1998) [hep-th/9804166].
- [45] E. A. Bergshoeff, O. Hohm, J. Rosseel, E. Sezgin and P. K. Townsend, “On critical massive (super)gravity in  $adS_3$ ,” *J. Phys. Conf. Ser.* **314**, 012009 (2011) [arXiv:1011.1153 [hep-th]].
- [46] S. M. Kuzenko and S. Theisen, “Correlation functions of conserved currents in  $N = 2$  superconformal theory,” *Class. Quant. Grav.* **17**, 665 (2000) [hep-th/9907107].
- [47] D. Butter and S. M. Kuzenko, “New higher-derivative couplings in 4D  $N = 2$  supergravity,” *JHEP* **1103**, 047 (2011) [arXiv:1012.5153 [hep-th]].
- [48] J. Louis and A. Micu, “Type II theories compactified on Calabi-Yau threefolds in the presence of background fluxes,” *Nucl. Phys. B* **635**, 395 (2002) [hep-th/0202168].
- [49] G. Dall’Agata, R. D’Auria, L. Sommovigo and S. Vaula, “ $D = 4$ ,  $N = 2$  gauged supergravity in the presence of tensor multiplets,” *Nucl. Phys. B* **682**, 243 (2004) [arXiv:hep-th/0312210].
- [50] R. D’Auria, L. Sommovigo and S. Vaula, “ $N = 2$  supergravity Lagrangian coupled to tensor multiplets with electric and magnetic fluxes,” *JHEP* **0411**, 028 (2004) [hep-th/0409097].
- [51] R. D’Auria and S. Ferrara, “Dyonic masses from conformal field strengths in  $D$  even dimensions,” *Phys. Lett. B* **606**, 211 (2005) [arXiv:hep-th/0410051].
- [52] J. Louis and W. Schulgin, “Massive tensor multiplets in  $N = 1$  supersymmetry,” *Fortsch. Phys.* **53**, 235 (2005) [arXiv:hep-th/0410149].
- [53] U. Theis, “Masses and dualities in extended Freedman-Townsend models,” *Phys. Lett. B* **609**, 402 (2005) [arXiv:hep-th/0412177].
- [54] S. M. Kuzenko, “On massive tensor multiplets,” *JHEP* **0501**, 041 (2005) [hep-th/0412190].
- [55] S. M. Kuzenko and J. Novak, “On curvature squared terms in  $N=2$  supergravity,” *Phys. Rev. D* **92**, no. 8, 085033 (2015) [arXiv:1507.04922 [hep-th]].